A Dynamic Interstellar Medium: 
Recent Numerical Simulations

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Abstract. Recent numerical simulations of the interstellar medium driven by energy input from supernovae and stellar winds indicate that H I clouds can be formed by compression in shock waves and colliding turbulent streams without any help from thermal and gravitational instabilities. The filling factor of the hot phase in these models does not exceed 20–40% at the midplane. The hot gas is involved in a systematic vertical outflow at |z| < 1–3 kpc, similar to that expected for galactic fountains, whereas the warm component may remain in hydrostatic equilibrium. The turbulent velocity is larger in the warmer phases, being 3, 10 and 40 km s$^{-1}$ in the cool, warm and hot phases, respectively, according to one of the simulations. The models exhibit global variability in the total kinetic and thermal energy and star formation rate at periods of (0.4–4) × 10$^8$ yr. Current models are still unable to reproduce dynamo action in the interstellar gas; we briefly discuss implications of the dynamo theory for turbulent interstellar magnetic fields.

1 Introduction

Recent numerical simulations of the interstellar medium (ISM) heated by supernovae and stellar winds have highlighted the importance of numerical experiment in the studies of the multiphase ISM. The simulations play a complimentary rôle to observational and semi-analytical studies and are especially important in clarifying such long-standing problems as the filling factors of the various phases of the ISM, the nature and lifetimes of H I clouds, dynamo action, etc. Simulations capturing relevant physical effects at appropriate resolution have become feasible only recently, and even first results reviewed here provide insight into many controversial issues.

The range of gas temperatures and densities in the ISM is tantalizingly wide, from $T \approx 10^6$ K, $n \approx 10^{-3}$ cm$^{-3}$ in the hot phase to 10 K and more than 10$^3$ cm$^{-3}$ in molecular clouds. The gas is involved in random motions over a wide range of scales (e.g., Spangler, 1999). It is still impossible to cover the whole range of parameter variation (i.e., all the phases) in a single model of the ISM. In this paper I shall review models of the diffuse phases of the ISM with densities below 10–100 cm$^{-3}$. Simulations of MHD turbulence in dense clouds have been discussed by Padoan et al. (1997, 1998), Stone et al. (1998), Ostriker et al. (1999) and Heitsch et al. (1999); an extensive review and references can be found in Vázquez-Semadeni et al. (2000).
2 Models of the multiphase ISM driven by supernovae

In this paper we discuss the results of three models of the ISM briefly summarized in Table 1. The models include similar physical effects and their results agree in many respects. All are non-linear, hydrodynamic or magnetohydrodynamic models that include external gravity (typical of the Solar neighbourhood), suitably idealized heating sources and radiative cooling (assuming partially ionized, optically thin gas). All models employ artificial viscosities to avoid excessively large gradients in the simulated quantities. Neither of the models includes explicitly ionization balance and cosmic rays.

| Dimension     | RBN\(^{(a)}\) 2D: \(xy\) and \(xz\) | VSP\(^{(b)}\) 2D: \(xy\) | KBSTN\(^{(c)}\) 3D: \(xyz\) |
|---------------|----------------------------------|--------------------------|---------------------------|
| Box size, kpc | \(2 \times 2\) and \(2 \times 15\) | \(1 \times 1\) \(0.5 \times 0.5 \times 2\) | \(1 \times 10\) \(2 \times 10\) |
| Resolution, pc| \(10\) and \(10 \times (10–1000)\) | \(2\) \(8\) | \(2\) \(8\) |
| Heating       | SN II, winds                      | Diffuse, winds           | SN II, SN I               |
| Temperature range, K | \(10^2–10^8\)                 | \(10^2–4 \times 10^4\)   | \(10^2–10^8\)            |
| Magnetic fields | –                                | imposed                 | self-excited             |
| Self-gravity  | –                                | yes                     | –                         |
| Rotation      | –                                | yes                     | yes                      |
| Cosmic Rays   | –                                | –                       | –                         |
| Ionization Balance | –                                | –                       | –                         |

References

\(^{(a)}\)Rosen et al. (1993, 1996), Rosen & Bregman (1995);
\(^{(b)}\)Vázquez-Semadeni et al. (1995), Passot et al. (1995), Ballesteros-Paredes et al. (1999);
\(^{(c)}\)Korpi et al. (1999a,b).

The model of Rosen et al. (1993, 1996) and Rosen & Bregman (1995) (labelled RBN) uses a finite-difference ZEUS code in two dimensions and neglects magnetic fields and rotation. These authors have results for a region in the Galactic plane (\(xy\)) and in a vertical plane (\(xz\)); the mesh size in the \(z\)-coordinate is variable, growing with \(z\) from 10pc to 1 kpc. A unique feature of this model is that the stellar Population I is included explicitly via equations for the stellar fluid coupled to the gas. The heating sources are Type II SNe, which inject energy instantaneously, and stellar winds modelled as a continuous heat source; both are correlated with the stellar density. To alleviate numerical problems, the simulations have been split into alternating stages, the interaction phase of the gas and stars when any motions are neglected, and a stage of hydrodynamic evolution when any coupling between stars and gas is neglected. SN explosions produce very hot gas, so the model contains three phases, i.e., the hot, warm and cool gas.

Vázquez-Semadeni et al. (1995, 1996, 1997, 1998), Passot et al. (1995), Scalo et al. (1998) and Ballesteros-Paredes et al. (1999), labelled VSPP in Table 1 (see also
earlier results in Passot et al., 1988, Léorat et al., 1990, and Vázquez-Semadeni & Gazol, 1995), employ a pseudo-spectral numerical scheme and restrict themselves to two dimensions in the Galactic plane. Their best runs have a spatial resolution of 2 pc or less, but only features at scales above $\sim 10$ pc are reliable (Passot et al., 1995). To avoid numerical difficulties, this model includes diffuse heating (which is thought to model the UV heating of the ISM) and stellar heating, but each stellar energy release event is spread over a period of $6 \times 10^{6}$ yr. Furthermore, both the heating and cooling rates are reduced by a factor 7–10 below the realistic values. As a result, only the warm and cool gas phases are generated, with the maximum gas temperature of $4 \times 10^{4}$ K. A peculiar feature of this model is the inclusion of self-gravity; however, its rôle is only marginal at the densities reached in the simulations. The model includes magnetic field, but any dynamo action is precluded in two dimensions.

The model of Korpi et al. (1999a,b) (KBSTN) is fully three-dimensional and includes the effects of rotation and magnetic field at a resolution of 8 pc. This is the only of the three models admitting realistic behaviour of vorticity and magnetic field including vortex stretching and dynamo action. The gas is heated by SNe modelled as thermal energy release over one time step which can be as short as 10–100 yr. As a result, the gas is heated to $T \simeq 10^{8}$ K at the explosion site and hot, warm and cool gas phases can be identified in the simulations. Both Type I and II SNe are included, which differ in both the occurrence rate and the vertical distribution. The SN explosions occur at randomly chosen sites, but Type I SNe can only explode in those regions where the gas density exceeds the average at that height. With this prescription, about 70% of the Type I SNe are clustered resembling OB associations. The simulations have only been run over time span of the order of $10^{8}$ yr, so it is yet unclear whether or not any magnetic fields can be supported by dynamo action in this model. A related model is presented by Gudiksen (1999) where an intricate algorithm for choosing SN explosion sites involving modelling of the stellar initial mass function has been developed.

Despite different numerical approaches employed and certain differences in the physical content, these models yield many coherent results which appear to be model-independent and which are our subject in what follows.

3 The multi-phase structure and cool gas filaments

In all the models, the cooling function is truncated at $T = 100$–300 K to avoid thermal instability at $T < 10^{5}$ K. Self-gravity is either neglected or unimportant. Nevertheless, cool, dense gas clouds with $T \lesssim 10^{3}$ K, $n = 1$–100 cm$^{-3}$ are a typical feature of all the simulations. The clouds are elongated and are better described as filaments or sheets (which can be distinguished from filaments only in 3D models); rounder structures occur at the intersections of the filaments. The filamentary H$\textsc{i}$ clouds occur at positions of converging turbulent flow, $\nabla \cdot \mathbf{v} < 0$ and are not gravitationally bound (Vázquez-Semadeni et al., 1995; Ballesteros-Paredes et al., 1999). This indicates that they are formed by compression (ram pressure) in the turbulent velocity field. The lifetime of the filaments is shorter than $10^{8}$ yr (Rosen & Bregman, 1995), which is close to the kinematic time scale. The filaments are mainly destroyed by mutual collisions.
In models with SNe, a more obvious mechanism of the cloud formation is compression in expanding shocks (Korpi et al., 1999). The length of the filaments is 50–100 pc and they are often extended vertically because of the vertical outflow of the hot gas, resembling H I ‘worms’ observed in spiral galaxies (Rosen & Bregman, 1995; Korpi et al., 1999). The three-dimensional model of Korpi et al. (1999) has both filaments and sheets of cool gas. In the model of Rosen & Bregman (1995), where SNe are confined to a rather thin layer, the cool filaments are not uncommon at $z = 1–3$ kpc, but in the model of Korpi et al. (1999) they are confined to $|z| \lesssim 100$ pc because Type II SNe that can occur at large heights prevent cool gas from rising higher.

As could be expected, the hot phase ($T > 10^5$ K) can only be generated by the SNe. The hot gas fills bubbles surrounded by dense shells; it breaks through the warm layer and streams into the halo (or beyond the upper boundary of the computational domain) at systematic velocities of 100–200 km s$^{-1}$ (Korpi et al., 1999); Rosen & Bregman (1995) obtain 400 km s$^{-1}$ at $z = 2–3$ kpc. The volume filling factors of the hot, warm and cold phases are 40–30%, 40–60% and 20–10% at $z = 1–3$ kpc in the model of Rosen & Bregman (1995). The filling factor of the hot gas increases at larger heights at the expense of the other phases. The filling factor of the hot phase in the model of Korpi et al. (1999) grows from 20–30% at $z = 0$ to 80–100% at $|z| = 1$ kpc; the remaining part of the volume is mostly occupied by the warm gas. Unlike the hot gas, the warm component is, on average, in hydrostatic equilibrium with a scale height of 200 pc; this agrees with a detailed analysis of hydrostatic equilibrium in the Solar vicinity by Fletcher & Shukurov (1999). As mentioned above, the cold gas is confined to $|z| \lesssim 100$ pc. Rosen & Bregman (1995) perform detailed fitting of the vertical gas distribution in different phases. The scale heights of the cold, warm and hot gas in one of their models (Run E) are 225, 550 and 2000 pc, respectively. Note that the gas called ‘cold’ in the model, can be identified with the cold and warm neutral medium, whereas the ‘warm’ gas corresponds to the warm ionized medium.

The phases of the modelled ISM are in a rough thermal pressure equilibrium. For example, density and temperature span four orders of magnitude in the model of Korpi et al. (1999), but thermal pressure varies in space only by an order of magnitude. Similarly, the density contrast is a factor of 50 whereas the pressure contrast is only a factor of 5 in the simulations of Vázquez-Semadeni et al. (1995). The approximate pressure equilibrium is not trivial to explain since the sound crossing time is not shorter than the dynamic time scale because both the turbulent and systematic motions are transonic or supersonic. Vázquez-Semadeni et al. (1995) discuss how the pressure balance can result from the instantaneous establishment of thermal balance in the gas. With their diffuse heating, both the cooling time and the heating time are 10–100 times shorter than the sound crossing time and the dynamic time scale, 10 pc/10 km s$^{-1} \simeq 10^7$ yr. Therefore, the heating rate $\Gamma$ must be always close to the cooling rate $\rho \Lambda$ (with $\rho$ the gas density); for $\Lambda \propto T^m$ (with $m$ a temperature-dependent index), the equilibrium temperature follows as $T_{eq} \propto (\Gamma/\rho)^{1/m}$. The resulting pressure (assuming ideal gas) is $P_{eq} \propto \Gamma^{1/m} \rho \gamma$, where $\gamma = 1–1/m = 0.33–0.66$ for $m = 1.5–2.9$. Since $\gamma < 1$, denser regions are cooler, and this results in weaker spatial variations of pressure (Vázquez-Semadeni, 1998; Vázquez-Semadeni et al., 2000). This explanation
relies on a non-localized nature of the heating (it also applies if the heating rate is weakly dependent on density—Passot et al., 1995), but it may be of more general importance.

4 Global variability

Quasi-periodic variations in the total thermal and kinetic energy (and also magnetic and gravitational energy where appropriate) accompanied by similar variations in star formation intensity are typical of the models. The period of the variations changes from one model to another in the range \((0.4–4) \times 10^8\) yr. Rosen & Bregman, 1995, do not discuss the global variability but mention oscillations, at a period \(2.5 \times 10^8\) yr, of the extent of the region occupied by filamentary gas structures. This is compatible with the vertical sound crossing time and free-fall time across 1 kpc, so these may be a global acoustic mode, although Vázquez-Semadeni et al. (1995) interpret the variations as an analogue of spiral density waves in their model.

5 Turbulence in the multiphase ISM

The models include a random element in the heating source, so it is not surprising that the gas motion is random, resembling developed turbulence. Vázquez-Semadeni et al. (1995, 1997) show that compressible motions with a kinetic energy spectrum \(E_k \propto k^{-2}\) characteristic of shock waves contribute significantly to the motions. Explosive heating by SNe driving shock waves would further increase this contribution. These authors argue that the velocity dispersion–size relation for interstellar clouds, \(\Delta v \propto R^{1/2}\) may result not from any kind of virial equilibrium, but rather directly from the turbulent spectrum: \(\nu_k^2 \propto k E_k \propto k^{-1}\), where \(\nu_k\) is velocity at a scale \(R \propto k^{-1}\).

The difference between cloud formation mechanisms in expanding SN shells and in colliding turbulent streams becomes significant when the vorticity of the motions, \(\omega\), is considered. Cool clouds formed by obliquely colliding streams of the turbulent flow must have enhanced vorticity: Vázquez-Semadeni et al. (1995) obtain an enhancement by a factor of 10 for denser clouds. Such an enhancement is not observed in a 3D model of Korpi et al. (1999a,b) where gas filaments and sheets can be also formed by expanding SN remnants and where essentially three-dimensional mechanisms are efficient in amplifying vorticity in all the phases. Vorticity is amplified at curved shock fronts distorted by ambient density inhomogeneities, by vortex stretching and by the baroclinic effect, and 60–90% of the turbulent energy is in vortical motions (Korpi et al., 1999a). Vortex stretching is an essentially three-dimensional effect; furthermore, the tendency for conservation of \(\omega/\rho\) in a two-dimensional flow also contributes to an artificial enhancement of vorticity in denser regions of a 2D model (Vázquez-Semadeni et al., 1995).

Korpi et al. (1999) computed the autocorrelation function of the vertical velocity for the cold, warm and hot phases separately. The total (three-dimensional) r.m.s. turbulent velocity of the cool gas is as low as \(3 \text{ km s}^{-1}\), whereas it is close to \(10 \text{ km s}^{-1}\) for the warm gas (approximately the speed of sound in the warm phase). The corre-
lation radius in the warm gas is about 30 pc (this is half the eddy size) and it only weakly changes with \( z \); this is interpreted as an indication of approximate vertical hydrostatic equilibrium of the warm gas with a scale height of about 200 pc. The r.m.s. random velocity in the hot gas, 40 km s\(^{-1}\) at \(|z| \ll 1\) kpc, is significantly larger than in the other phases, but still remains significantly smaller than the sound speed in the hot phase, 100 km s\(^{-1}\). The vertical streaming of the hot gas discussed above can feed turbulence higher in the halo, so the turbulent velocity of 60 km s\(^{-1}\) observed in the halo by Kalberla et al. (1998; see also Pietz et al., 1998) seems to be perfectly compatible with the model. (Turbulent velocities of order 100 km s\(^{-1}\) in the halo have been invoked in galactic dynamo models involving the halo—Sokoloff & Shukurov, 1990; Brandenburg et al., 1992, 1993.) The typical horizontal radius of the hot regions at the midplane is about 20 pc; this can be identified with a chimney radius. The correlation radius for the hot phase grows with \( z \) together with the volume filling factor.

6 Magnetic fields and dynamo action

The only model admitting hydromagnetic dynamo action among the three models discussed is that of Korpi et al. (1999a,b). These simulations are initialized with a weak (0.1 \( \mu \)G) regular magnetic field to provide a seed field for the dynamo. Both the large-scale and the random magnetic fields show a tendency to grow in the model, but runs spanning at least \( 10^9 \) yr are required to demonstrate any dynamo action convincingly.

Magnetic fields in the 2D simulations of Passot et al. (1995) do not decay only because the model has an imposed uniform magnetic field, maintained throughout the simulations, which is tangled by small-scale motions. Magnetic fields in the cool filaments formed by converging flows often have a reversal along the filament axis. Magnetic energy in these simulations is close to the kinetic energy in solenoidal motions at all scales. Magnetic fields, like vorticity, are especially sensitive to the dimensionality of the model, and two-dimensional results should be considered with caution.

According to dynamo theory, the growth time of the random magnetic field due to the small-scale dynamo action is expected to be as short as the eddy turnover time, \( \simeq 10^7 \) yr in the warm phase. A plausible statistically steady state resulting from the saturation of the dynamo action is an ensemble of magnetic flux ropes (e.g., Zeldovich et al., 1983, Sect. 8.IV). The length of the rope is of the order of the turbulent scale, \( l \simeq 50–100 \) pc, and the thickness is of the order of \( lR_m^{-1/2} \), where \( R_m \) is the magnetic Reynolds number. Dynamo action can occur provided \( R_m > R_{m,cr} \), where the critical magnetic Reynolds number is estimated as \( R_{m,cr} = 60–100 \) in simplified models of homogeneous, incompressible turbulence.

Subramanian (1999) suggested that a steady state, reached via the back-action of the magnetic field on the flow, can be established by the reduction of the effective magnetic Reynolds number \( R_{m,eff} \) down to the value critical for the dynamo action.

\(^1\)In this context, ‘small’ scales are understood as those smaller than the correlation scale of the turbulence, 50–100 pc in the warm gas.
Then $R_{m,\text{eff}} \approx R_{m,\text{cr}}$ in the steady state. Therefore, the thickness of the ropes in the steady state can be estimated as $l R_{m,\text{cr}}^{-1/2}$. Using a model nonlinearity in the induction equation with incompressible velocity field, Subramanian (1999) showed that the magnetic field strength within the ropes $b_0$ saturates at the equipartition level with kinetic energy density, $b_0^2/8\pi \simeq \frac{1}{2} \rho v^2$, where $\rho$ is the gas density and $v$ is the r.m.s. turbulent velocity. The average magnetic energy density is estimated as $\langle b^2/8\pi \rangle \simeq C R_{m,\text{cr}}^{-1} \frac{1}{2} \rho v^2$, where $C$ is a numerical coefficient, presumably of order unity. The volume filling factor of the ropes then results as $f \simeq R_{m,\text{cr}}^{-1} \simeq 0.01$; correspondingly, the mean magnetic energy generated by the small-scale dynamo in the steady state is about 1% of the turbulent kinetic energy density.

Using parameters typical of the warm phase of the ISM, this theory predicts that the small-scale dynamo would produce magnetic flux ropes of the length (or the curvature radius) of about $l \simeq 50$–100 pc and thickness $l R_{m,\text{cr}}^{-1/2} \simeq 5$–10 pc. The field strength within the ropes, if at equipartition with the turbulent energy, has to be of order $1.5 \mu G$ in the warm phase ($n = 0.1 \text{ cm}^{-3}$, $v = 10 \text{ km s}^{-1}$) and $0.5 \mu G$ in the hot gas ($n = 10^{-3} \text{ cm}^{-3}$, $v = 40 \text{ km s}^{-1}$). Note that other models of the small-scale dynamo admit solutions with magnetic field strength within the ropes being significantly above the equipartition level (Belyanin et al., 1993).

The small-scale dynamo is not the only mechanism producing turbulent magnetic fields (e.g., Beck et al., 1996, Sect. 4.1 and references therein). Any mean-field dynamo action producing magnetic fields at scales exceeding the turbulent scale also generates small-scale, turbulent magnetic fields. Similarly to the mean magnetic field, this component of the turbulent field presumably has a filling factor close to unity in the warm gas and its strength is expected to be close to equipartition with the turbulent energy at all scales. (As argued below, this component of the turbulent magnetic field may be confined to the warm gas, so magnetic field in the hot phase may have a better pronounced ropy structure.)

The overall structure of the interstellar turbulent magnetic field in the warm gas can be envisaged as a quasi-uniform fluctuating background with one percent of the volume occupied by flux ropes (filaments) of a length 50–100 pc with well-ordered magnetic field. The ropes can plausibly produce elongated, localized filamentary structures of polarized emission and Faraday rotation recently observed in the Milky Way by Wieringa et al. (1993) and Gray et al. (1998). The basic distribution described would be further complicated by compressibility, shock waves, MHD instabilities (such as Parker instability), etc.

The site of the mean-field dynamo action is plausibly the warm phase rather than the other phases of the ISM. The warm gas has a large filling factor (so it can occupy a simply connected global region), it is, on average, in a state of hydrostatic equilibrium, so it is an ideal site for both the small-scale and mean-field dynamo action. Molecular clouds and dense H I clouds have too small a filling factor to be of global importance. Fletcher & Shukurov (1999) argue that, globally, molecular clouds can be only weakly coupled to the magnetic field in the diffuse gas. (However, the small-scale dynamo can work within the clouds.) The time scale of the small-scale dynamo in the hot phase is $\gtrsim 1/v \simeq 10^6 \text{ yr}$ for $l = 40 \text{ pc}$ and $v = 40 \text{ km s}^{-1}$. This can be shorter
than the advection time due to the vertical streaming, $h/V_z \simeq 10^7$ yr with $h = 1$ kpc and $V_z = 100$ km s$^{-1}$. Therefore, the small-scale dynamo action should be possible in the hot gas even at small heights. However, the growth time of the mean magnetic field must be significantly longer than $l/v$ reaching a few hundred Myr (e.g., Beck et al., 1996, Sect. 4.4). Thus, the hot gas can hardly contribute significantly to the the mean-field dynamo action in the disc and can drive the dynamo only in the halo. The fountain flow rather pumps the mean magnetic field out from the disc (Brandenburg et al., 1995); this effect can contribute to the saturation of the mean-field dynamo in the disc.

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