LARGE-SCALE EXPLOSIONS AND SUPERBUBBLES IN THE GALACTIC DISK AND HALO. I. MAGNETOHYDRODYNAMIC SIMULATIONS

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ABSTRACT

The evolution of the interstellar superbubbles arising from sequential supernova explosions or winds from OB associations in the Galactic disk is studied numerically, using a two-dimensional magnetohydrodynamics (MHD) code. We find that in the presence of horizontal magnetic fields of strength comparable to that in the Galactic disk, $B \approx 5 \mu G$, the vertical expansion of the superbubbles (the contact surface) can, under some conditions, be significantly inhibited by the effect of a decelerating $J \times B$ force. At the same time, the outermost effect of the disturbance actually propagates somewhat faster than in nonmagnetic cases, as an MHD fast shock or nonlinear wave. The implications of our results for galactic supershells, the galactic fountain, observed activity in starburst galaxies, and supernova remnants are briefly discussed.

Subject headings: ISM: bubbles — ISM: magnetic fields — MHD — supernova remnants

I. INTRODUCTION

It is now believed that some of the observed large-scale expanding superstructures of interstellar matter, called superbubbles (see Tenorio-Tagle & Bodenheimer 1988 for a review), such as those in the galactic disks of our galaxy, M31, and the Large Magellanic Cloud, are formed by multiple supernovae and winds from OB associations. Superbubbles are filled with hot ($\sim 2 \times 10^6$ K) coronal plasmas emitting soft X-rays (e.g., Hayakawa 1979), and surrounded by dense cold shells, called supershells, emitting the H I 21 cm line (Heiles 1984). Their sizes range from 300 pc to 1 kpc, and the total energies required to explain them are estimated to range from $5 \times 10^{54}$ to $10^{54}$ ergs (Tenorio-Tagle & Bodenheimer 1988). As such, they are too energetic to be explained by one single supernova explosion ($E \sim 10^{51}$ ergs).

Superbubbles are thought to play a role in supplying hot plasma (including both mass and energy) to the Galactic halo when they blow out of the disk into the halo. If such hot plasma can cool radiatively and thereby form dense clouds in the halo, these clouds would fall back toward the disk and might explain the well-known high-velocity clouds observed in 21 cm emission. This cyclical process involving the injection of hot gas into the halo and the return of halo gas to the disk following radiative cooling and cloud formation has been called the “Galactic fountain” (Shapiro & Field 1976).

Superbubbles are suggested to exist in other galaxies, as well. Aside from their role in normal disk galaxies, they may be responsible for some of the observed activity of starburst galaxies, such as M82 (e.g., Nakai et al. 1987; Sofue 1988; Umemura et al. 1988).

Tomisaka & Ikeuchi (1986) performed two-dimensional hydrodynamic simulations of superbubbles in the Galactic disk and found that in the case of low interstellar gas density in the Galactic central plane, $n_0 \sim 0.1$ cm$^{-3}$, superbubbles blow out of the disk, into the halo. Ikeuchi (1986) referred to these vertically elongated superbubbles as “chimney” because, aside from their elongation, these bubbles channel “dirty” (full of heavy elements), hot plasma into the halo. Tenorio-Tagle, Bodenheimer, & Rózycka (1987) and Mac Low, McCray, & Norman (1989) have made similar simulations but with higher resolution and accuracy. On the basis of their results, Mac Low, McCray, & Norman (1989) suggested that the low Galactic halo actually consists of a froth of merged superbubbles. Mac Low & McCray (1988) and Bisnovati-Kogan, Blinnikov, & Silich (1989) discussed the evolution of superbubbles in terms of simpler, approximate methods which involved a numerical extension of the analytical approximation for adiabatic point explosions in an exponential atmosphere by Kompaneets (1960). They claimed that these approximate methods can reproduce the numerical simulations quite well.

Although there have been several studies of the dynamical evolution of superbubbles, the effect of magnetic fields is not fully understood yet. Recent observations, however, suggest the presence of large-scale (mainly horizontal) magnetic fields with average field strength $B \sim 3-8 \mu G$ in the disks of many spiral galaxies (see, e.g., Sofue, Fujimoto, & Wielebinsky 1986; Ruzmaikin, Shukurov, & Sokoloff 1987; Rand & Kulkarni 1989). The field strength is even stronger in the central regions of galaxies, say 10–100 $\mu G$ (Sofue & Fujimoto 1989). On average, therefore, the magnetic pressure is comparable to the thermal gas pressure in the interstellar medium and may even dominate in some regions (Cox 1988): the gas pressure is $p_{\text{gas}} \sim 10^{-12} n_0 T_4$ for density $n = n_0$ cm$^{-3}$ and temperatures $T = 10^4 T_4$ K, whereas the magnetic pressure is $p_{\text{mag}} \sim 10^{-12} B_4^2$ for magnetic field $B = 5 B_4 \mu G$. In any case, one cannot generally neglect the presence of the magnetic fields, but should instead expect various kinds of new phenomena which are not known in pure hydrodynamic situations: e.g., Parker instability (Parker 1966; Mouschovias 1974; Horiuchi et al. 1988; Matsumoto et al. 1988), and MHD collimation of astrophysical jets (Shibata & Uchida 1985).

The evolution of superbubbles may also be influenced by the magnetic field (Bernstein & Kulsrud 1965; Kulsrud et al. 1965; Mineshige & Shibata 1990). Consider an explosion or sequence of explosions which deposits an energy $E_{\exp}$ in a gas with magnetic...
field strength $B$. Very crudely, we can expect the magnetic field to be dynamically important if the superbubble reaches a size comparable to the equipartition radius, $r_e$, given by

$$E_{\text{exp}} = \frac{B^2}{8\pi} \frac{4\pi}{3} r_e^3 \quad (1.1)$$

(if there is no radiation.) In terms of $E_{53} = E_{\text{exp}}/(10^{53} \text{ ergs})$, this gives

$$r_e \approx 1E_{53}^{1/3} B_0^{-2/3} \text{ kpc} \quad (1.2)$$

This is a length scale of relevance to superbubbles which expand into the galactic halo.

Umemura et al. (1988) have provided an example of the dynamical effect of magnetic fields in their simulations of an explosion in the Galactic center, assuming that the magnetic fields there were initially uniform and perpendicular to the Galactic plane. They found that, for an explosion energy $E = 10^{54} \text{ ergs}$ and a field strength of $B \sim 30–100 \mu\text{G}$, the flow forms a hollow, cylindrical shell structure, similar to radio lobes observed in our galactic center and in M82. For even higher field strength, the mass motion was found to be highly collimated, producing a jet with a narrow opening angle along vertical field lines.

Away from the Galactic center, however, the magnetic fields are thought to be, on average, horizontal to the Galactic plane. Hence, we examine in this paper the effects of such horizontal magnetic fields on the evolution of superbubbles, using two-dimensional magnetohydrodynamic (MHD) numerical simulations. We will see significant suppression of the expansion of a superbubble toward the halo because of the decelerating $J \times B$ force even in the adiabatic cases (see Tomisaka 1990); break-out of the disk, such as has been described in terms of the formation of a “chimney” can, therefore, be prevented in the presence of the large-scale horizontal magnetic fields.

Preliminary results of these calculations were previously summarized in Shapiro (1990) and Shapiro, Mineshige, & Shibata (1990). Our calculations here are of higher numerical resolution but agree qualitatively with those earlier results. In the meantime, in an independent effort, Tomisaka (1990) has treated this problem numerically in three dimensions. Our two-dimensional calculations with $241 \times 361$ grid cells is of higher resolution (per dimension) than Tomisaka's three-dimensional calculations of $81 \times 81 \times 81$ cells.

In the subsequent paper (Shapiro, Mineshige, & Shibata 1993) we will construct a simple analytical model based on the Kompaneets approximation, and discuss the conditions necessary for magnetic fields to be dynamically important in relation to the present work.

The plan of this paper is as follows: we give the physical assumptions, basic equations, and numerical procedures in § 2. The results of the numerical simulations both of nonmagnetic and magnetic cases are shown in § 3. We will see appreciable differences between the nonmagnetic and magnetic models. Section 4 is devoted to discussion and summary, while the implications of these results for the Galactic fountain, supershells, superbubbles in the Galactic center and starburst galaxies, and supernova remnants are briefly discussed in § 5.

2. BASIC EQUATIONS AND NUMERICAL METHODS

2.2. Assumptions and Basic Equations

Our calculations are limited to two-dimensional. As such, in order to describe an explosion in the Galactic disk, with horizontal magnetic field along the local azimuthal ($\theta$) direction in galactocentric ($r, \theta, z$) cylindrical coordinates, we adopted Cartesian ($x, y, z$) coordinates, where $\hat{x} = \hat{\theta}$ and $\hat{y} = -\hat{r}$ for some fixed value of $\theta$. In that case, $\mathbf{B} = B_0 \hat{x}$ is locally tangential in the disk geometry (see Fig. 1).

We make the following assumptions in this paper: (1) the medium is an ideal gas with ratio of specific heat $\gamma = 5/3$; (2) the magnetic field is frozen into the gas; (3) the gravitational acceleration is either constant or proportional to $z > 0$; (4) the galactic differential rotation can be ignored; (5) only two-dimensional motions are allowed; (6) the radiative cooling and thermal conduction are ignored. Cooling time-scale is roughly $\tau_{\text{cool}} \sim 3 \times 10^5 (T/10^8 \text{ K})^{1/2} (n/10^{-4} \text{ cm}^{-3})^{-1} \text{ Myr}$ for hot interiors and $\tau_{\text{cool}} \sim 1.0 (T/10^5 \text{ K})^{1/2} (n/1 \text{ cm}^{-3})^{-1} \text{ Myr}$ for compressed shells, where $1 \text{ Myr} = 10^6 \text{ yr}$ and we used the bremsstrahlung cooling. This is hence usually much longer than those of model calculations ($\sim$ a few Myr) presented in this paper except at cold shells (will be discussed later in § 3.3).

By assumption (5), $y$ components of the velocity and the magnetic field, $v_y, B_y$, and the derivatives with respect to $y$ are taken to be zero. The basic equations are then of the mass conservation, given by

$$\frac{\partial \rho}{\partial t} + \frac{\partial}{\partial x} (\rho v_x) + \frac{\partial}{\partial z} (\rho v_z) = 0 \quad (2.1)$$

Fig. 1.—Adopted coordinates
momentum conservation, given by

$$\frac{\partial}{\partial t} (\rho v_x) + \frac{\partial}{\partial x} \left[ \rho v_x^2 \frac{p}{\gamma - 1} + \frac{1}{2} \rho (v_x^2 + v_z^2) + \frac{1}{8\pi} (B_x^2 + B_z^2) \right] + \frac{\partial}{\partial z} \left( \rho v_x v_z - \frac{1}{4\pi} B_x B_z \right) = 0,$$

(2.2)

Faraday's law in the MHD approximation, given by

$$\frac{\partial}{\partial t} B_x + \frac{\partial}{\partial z} (v_x B_x - v_z B_z) = 0,$$

(2.4)

and energy conservation, given by

$$\frac{\partial}{\partial t} \left[ \frac{p}{\gamma - 1} + \frac{1}{2} \rho (v_x^2 + v_z^2) + \frac{1}{8\pi} (B_x^2 + B_z^2) \right] + \frac{\partial}{\partial x} \left[ \frac{\gamma}{\gamma - 1} \rho v_x \left( v_x^2 + v_z^2 \right) + \frac{1}{4\pi} \left( v_x B_x - v_z B_z \right)^2 \right] + \frac{\partial}{\partial z} \left[ \frac{\gamma}{\gamma - 1} \rho v_z \left( v_x^2 + v_z^2 \right) + \frac{1}{4\pi} \left( v_x B_x - v_z B_z \right)^2 \right] + \rho g v_z = 0,$$

(2.6)

where $g$ is the gravitational acceleration, $p$ is the gas pressure, and other symbols have their usual meaning.

2.2. Initial Conditions

We consider a gas layer which is initially uniform in the $x$-direction. For the initial vertical density and temperature distributions, we adopt the following different models.

2.2.1. Exponential Disk, with and without Field

In the isothermal exponential model (hereafter “exponential model” or model E), temperature is assumed to be constant, $T = 10^4$ K. The magnetic field is initially horizontal, and the plasma $\beta$ defined by $\beta = p_{\text{gas}}/p_{\text{mag}}$, where $p_{\text{gas}}$ is the gas pressure given below and $p_{\text{mag}} = B^2/8\pi$, is kept constant. So we have

$$B_x(z) = \left[ \frac{8\pi p_{\text{gas}}(z)}{\beta} \right]^{1/2}.$$

(2.7)

In this case, since $p_{\text{gas}}$ is a function of height, the initial magnetic field is not uniform.

We assume constant gravity in the direction perpendicular to the disk plane. The initial density distributions are then obtained by integrating the magnetohydrostatic balance equation,

$$\frac{\partial}{\partial z} \left[ \rho(z) T(z) \left( 1 + \frac{1}{\beta} \right) \right] + \rho(z) g = 0.$$

(2.8)

The density has therefore an exponential profile:

$$\rho(z) = \rho_0 \exp \left( -z/H_0 \right),$$

(2.9)

where $n_0 \equiv \rho_0/m_p \approx 10^3$ cm$^{-3}$ ($m_p$ denotes the proton mass) is the number density at $z = 0$, and the scale height of the disk $H_0 = 100$ pc. The strength of horizontal magnetic field hence decreases with $z$ as

$$B_x(z) = B_0 \exp \left( -z/2H_0 \right),$$

(2.10)

with $B_0 = (8\pi \rho_0 T_0/\mu B)^{1/2} \approx 4.5 \beta^{-1/2} \mu G$. Finally, the gas pressure is given by

$$p_{\text{gas}}(z) = \rho(z) \mathcal{R} T(z).$$

(2.11)

Here $\mathcal{R}$ is the gas constant, and $\mu$ is the mean molecular weight.

This value for the scale height is close to the effective thickness of the Galactic disk H I cloud layer as described by Spitzer (1978) and is chosen to facilitate comparisons of our magnetized results with previous nonmagnetic calculations such as those of Mac Low, McCray, & Norman (1989) and Tomisaka & Ikeuchi (1986). It has also been suggested that observations favor a hybrid atmosphere with two components, modeled as a Gaussian with scale height less than 200 pc plus an exponential with a scale height of 500 pc (see, e.g., Lockmann, Hobbs, & Shull 1986). Since our results are intended to illustrate the important new effects of the inclusion of magnetic fields, we shall restrict attention here to the simpler one-component model and leave consideration of more complicated initial conditions like that of the hybrid model for future work.
2.2.2. Gaussian Disk, with and without Field

In the second model, we assume the linear dependence of gravity on \( z \), giving rise to a Gaussian density distribution in magnetohydrostatic equilibrium (hereafter "Gaussian model" or model \( G \)),

\[
\rho(z) = \rho_0 \exp \left( -\frac{z^2}{2H_0^2} \right),
\]

(2.12)

with \( \rho_0 = 1.0 m_p \) and \( H_0 = 100 \) pc. Again the temperature is assumed to be constant in space, \( T = 10^4 \) K. As with the exponential model, we calculate a nonmagnetic model and a magnetic model in which the magnetic field is initially parallel to the Galactic plane. Again the plasma \( \beta \) is assumed to be constant vertically, thus the field strength varies with height, \( z \), according to

\[
B_x(z) = B_0 \exp \left( -\frac{z^2}{4H_0^2} \right),
\]

(2.13)

with \( B_0 \approx 4.5 \) \( \mu \)G, so that the plasma \( \beta = 1 \).

2.2.3. Magnetostatic Disk Halo

Finally, we consider a model with a halo (hereafter "disk halo model" or model \( H \)). In this case, the pressure in the disk has a finite value and balances gravity. In order to solve the initial equilibrium structure, we first prescribe the temperature distribution:

\[
T(z) = T_{\text{disk}} + (T_{\text{halo}} - T_{\text{disk}})H\left(\frac{z - z_{\text{halo}}}{w_{\text{tr}}}\right),
\]

(2.14)

with the temperature of the disk \( T_{\text{disk}} = 10^4 \) K, the temperature of the halo \( T_{\text{halo}} = 10^5 \) K, the height of the disk-halo boundary \( z_{\text{halo}} = 100 \) pc, the width of the transition layer \( w_{\text{tr}} = 20 \) pc, and \( H(x) = \frac{1}{2} \left( 1 + \tanh \left( \frac{x}{x_0} \right) \right) \). The density profile is obtained by the magnetohydrostatic balance equation (2.8), and so within the disk, \( z < z_{\text{halo}} \), the density has an exponential profile (2.9) with the scale height \( H_0 = 100 \) pc, and abruptly decreases by a factor \( \sim 100 \) beyond the disk halo interface.

2.3. SN Explosion

Following Tomisaka & Ikeuchi (1986) we assume successive mass and energy injection by sequential supernova explosions. The mass and energy input per explosion are \( M_{\text{exp}} = 10 M_\odot \) and \( E_{\text{exp}} = 10^{51} \) ergs, and the time interval between injection is \( t_{\text{exp}} = 3.0 \times 10^5 \) yr. This yields a mechanical luminosity \( L_{\text{exp}} \approx 10^{38} \) ergs s\(^{-1}\).

Numerically we distribute the mass and energy for each explosion uniformly throughout a region with volume \( V_{\text{exp}} = \pi r_{\text{exp}}^2 y_{\text{exp}} \), where the explosion radius is \( r_{\text{exp}} = (x_{\text{exp}}^2 + z_{\text{exp}}^2)^{1/2} = 20 \) pc and \( y_{\text{exp}} = 500 \) pc. [Note that, since our model is essentially two-dimensional in order to study the effects of the horizontal magnetic field, the explosion is actually confined in the two-dimensional \((x, z)\) plane.] The density and the pressure enhancement for each of the successive explosions within the radius \( r = (x^2 + z^2)^{1/2} \leq r_{\text{exp}} \) are thus \( \rho_{\text{exp}} = (\gamma - 1)E_{\text{exp}}/V_{\text{exp}} \) and \( \rho_{\text{exp}} = M_{\text{exp}}/V_{\text{exp}} \), respectively.

2.4. Normalization

Equations (2.1)–(2.6) are nondimensionalized in the following way:

\[
(x, z) \rightarrow (x/H_0, z/H_0),
\]

(2.15a)

\[
(v_x, v_y) \rightarrow (v_x/c_{s,0}, v_y/c_{s,0}),
\]

(2.15b)

\[
\rho \rightarrow \rho / \rho_0,
\]

(2.15c)

\[
p \rightarrow p / (\rho_0 c_{s,0}^2),
\]

(2.15d)

\[
(B_x, B_z) = \left[ B_x / (\rho_0 c_{s,0}^2)^{1/2}, B_z / (\rho_0 c_{s,0}^2)^{1/2} \right],
\]

(2.15e)

and

\[
g \rightarrow g / (c_{s,0}^2/H_0),
\]

(2.15f)

with again \( H_0 = 100 \) pc, \( \rho_0 = 10^{-24} \) g cm\(^{-3}\), and \( c_{s,0} = 12 \) km s\(^{-1}\). We note that if we define \( \rho_0 = \rho_0 \Theta T_{\text{disk}}/\mu \), then \( \rho_0 = \rho_0 c_{s,0}^2/\gamma \).

2.5. Boundary Conditions and Numerical Procedures

We assume symmetry with respect to reflection about \( x = 0 \) and \( z = 0 \), and free boundaries for \( x = X_{\text{max}} \) and \( z = Z_{\text{max}} \). By definition, the fluid and waves should be able to pass freely through the free boundaries. It is, however, very difficult to design an ideal free boundary especially for the case of two- (or three)-dimensional problems (see e.g., Roche 1972); it is often observed that waves reflected at the free boundary crucially influence the dynamics occurring in the computing box, or that numerical errors generated at the free boundary lead to numerical instabilities. We here adopt first-order extrapolation for the free boundary (Chu & Sereny 1974; Shibata 1983):

\[
\frac{\partial \delta q}{\partial x} = 0
\]

(2.16a)

for the right boundary, and

\[
\frac{\partial \delta q}{\partial z} = 0
\]

(2.16b)
for the upper boundary, respectively, with $q$ being physical quantities, such as $\rho, p, v_x, v_y, B_x, B_y,$ etc., and $\delta q \equiv q(t + \delta t) - q(t).$ This method has the advantage that it is the simplest and is even better than other higher order schemes in some cases. Furthermore, this method is excellent for problems involving a strong shock wave or a supersonic (or super-Alfvénic) outflow, because no physical information will be transported back into the computing box from the boundary in these problems. Even in the case where the outflow is subsonic (or sub-Alfvénic) locally, this method gives acceptable results, as long as the plasma $\beta$ is greater than 0.01. In fact, Shibata & Uchida (1985, 1990) and Umemura et al. (1988) checked the effect of free boundaries in the calculations of MHD jets and blast wave shocks by changing the size of the computing boxes, and found no distinct differences there.

Equations (2.1)–(2.6) are then solved numerically by using a modified Lax-Wendroff scheme (Rubin & Burstein 1967) with artificial viscosity both in the $x$- and $z$-directions (Richtmyer & Morton 1967). The mesh spacings are $\Delta x = \Delta z = 2.5$ pc ($= H_0 /40$). The total number of mesh points is $(N_x \times N_z) = (241 \times 361).$ The total area is $(X_{\max} \times Z_{\max}) = (6H_0 \times 9H_0) = (600$ pc $\times 900$ pc).

The MHD code was tested in the following ways:

1. We first calculate successive supernova explosions in a uniform nonmagnetic medium using our code, and compare the results with two-dimensional Sedov's solution: $r_s \approx (E/\rho)^{1/4} t^{1/2}.$ The agreement is quite good, and numerical errors are less than 1%.

2. We then evaluate the values of $\text{div}(B)$ in the course of calculations, finding that $\text{div}(B)$ is less than 0.1 percent of $B_0 / \Delta x (= B_0 / \Delta z)$ except at the positions of shock waves where $\text{div}(B)$ amounts to $\sim 0.3$ percent of $B_0 / \Delta x$ at maximum.

3. Finally, we calculate the total mass and total energy in the calculated zone together with their inputs and outputs, confirming that both of these are conserved with accuracy less than 1% of their initial values.

Further description of the code and tests are in Shibata (1983).

3. NUMERICAL SIMULATIONS OF MAGNETIC SUPERBUBBLES

In this section, we give the results of both the nonmagnetic hydrodynamical calculations, which are similar to previous calculations (see Tomisaka & Ikeuchi 1986; Tenorio-Tagle & Bodenheimer 1988; Mac Low, McCray, & Norman 1989), and of the MHD simulations, in order to see the differences arising from the introduction of horizontal magnetic fields.

3.1. Exponential Model with No Magnetic Field: Model E0

Figures 2a and 2b display sequences of structural evolution of superbubbles for the exponential nonmagnetic model at two different times, 5.1 (Fig. 2a) and 10.2 Myr (Fig. 2b). Here 1 Myr $= 10^8$ yr. Contours of density are shown in the left panels of each figure by solid lines in each figure, with logarithmic spacing of 0.5 dex from $\rho = 10^{-28}$ to $10^{-24}$ g cm$^{-3}$. Gray zones indicate the positions of the forward shock front (outer ring) and the contact discontinuity (inner rings). Velocity vectors are also illustrated in the right panels by arrows, whose length are taken to be proportional to their absolute values. The norm vector, placed in the lower right edge of the right panel, corresponds to the upward velocity of $10 c_{s,0} \approx 120$ km s$^{-1}$.

Figure 3 depicts the vertical distributions at $x = 0$ of the pressure (upper left), the temperature (lower left), the $z$-component of the velocity divided by the local sound velocity (upper right), and the density (lower right) at times of 0.0, 5.1, and 10.2 Myr. A dense shell is formed between the contact discontinuity and the forward shock: a cold ($T \approx 10^5$ K) dense ($\rho \approx 10^{-24}$ g cm$^{-3}$) shell surrounds the hot ($T \approx 10^6$ K) tenuous ($\rho < 10^{-27}$ g cm$^{-3}$) gas (see, e.g., Tomisaka & Ikeuchi 1986). In Mac Low, McCray, & Norman (1989) a reverse shock is formed around the exploding area (see their Fig. 4). There is no clear reverse shock in our simulations, however. This is because we input mass and energy intermittently rather than continuously, and so is the reverse shock very transient; shortly after each explosion it propagates backward to the origin and dies away. In the model of Mac Low, McCray, & Norman, on the other hand, the continuous injection of mass at the origin makes the reverse shock persist.

3.2. Exponential Model with Horizontal Magnetic Field: Model E1

As anticipated by ours in § 1, the effects of magnetic fields on the evolution of superbubbles are appreciable when the superbubble diameter reaches a few hundred parsec for a horizontal, constant-$\beta$, preshock magnetic field of strength $B \approx 4.5 \mu G$ at $z = 0$. In Figure 4 we plot the evolution of the magnetic superbubble. The model parameters are the same as those in Figure 2 except that, for the case in Figure 4, there is horizontal magnetic fields with $\beta = 1$ initially. Magnetic field lines are displayed in the right panels. To have the same scale height, therefore, the strength of gravity ($g$) is taken twice as much as that in the nonmagnetic case, because

$$H_0 = \frac{\mathcal{R} T_0 [1 + (1/\beta)]}{\mu g},$$

(3.1)

from eq. (2.8). We also illustrate the sectional views of the pressure, the density, the $z$-component of the velocity, and the temperature along the z-axis in Figure 5. Contribution of magnetic pressure is depicted by the dashed lines in the upper left panel. We see $p_{\text{gas}} \approx p_{\text{mag}}$ in the cold dense shell.

We note here two main differences from the nonmagnetized model: acceleration of the shock front and broadening and deceleration of the contact surface. In Figure 6 we plot the heights of the shock front (dashed lines) and the contact discontinuity (solid lines) for models E0 and E1. We take two values for the field strengths: $1/\beta = 0$ and 1. The deceleration of the contact discontinuity (CD) and the acceleration of the shock front (SF) relative to the velocity in the nonmagnetized case are clear in this figure.

The acceleration of the shock can be explained as follows. Consider the propagation of the forward shock upward along the z-axis. This is a shock propagating perpendicular to the preshock $B = (B_x, 0, 0)$. Let the density jump be $D \equiv \rho_2 / \rho_1$, the (gas) pressure jump be $P \equiv p_2 / p_1$, let $B_1$ be the preshock value of $\beta$, and let $M_{1,1}$ be the isothermal Mach number of the shock relative to the preshock isothermal sound speed [i.e. $M_{1,1}^2 \equiv v_1^2/(p_1/\rho_1)$]. Using the MHD shock jump conditions, it can be shown that, in the
Fig. 2.—Evolution of superbubbles in the nonmagnetized medium with an exponential density profile (model E0). For model parameters, see Table 1. Contours of density are indicated in the left panels by thick lines with logarithmic spacing of 0.5 dex from $\rho = 10^{-25}$ to $10^{-28} \, \text{g cm}^{-3}$, and the velocity vectors are illustrated in the right panels by arrows. Gray zones in the left panel indicate the positions of the forward shock front (outer ring) and the contact discontinuity (inner rings). The norm vector (corresponding to the velocity of $100_\text{km s}^{-1}$) is also indicated in the right panel. Elapsed times are 5.1, and 10.2 (Myr: 1 Myr = 10^6 yr) for (a) and (b), respectively.

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limit of large shock strength (i.e., \( M_{t,1} \to \infty \), \( D \to (\gamma + 1)/(\gamma - 1) \) and \( P \to 2M_{t,1}^2/(\gamma + 1) \), the same as the nonmagnetic, strong shock jump conditions. However, for moderate values of \( M_{t,1} \), it can be shown, the presence of the magnetic field reduces the shock compression factor \( D \) and increases the shock velocity for a fixed \( P \). A comparison of the pressure profiles in Figures 3 and 5, however, suggests that the quantity which remains roughly the same for the two superbubble simulations, with and without a magnetic field, is not \( P \) but is instead \( P_{\text{tot,2}}/P_1 \equiv (p_2 + B_2^2/8\pi)/p_1 = P + (D^2/\beta_i) \). Nevertheless, we still expect the shock to accelerate relative to the unmagnetized case by the following argument. First, for a fixed value of \( P_{\text{tot,2}}/P_1 \), the postshock velocity of the undisturbed external medium is actually reduced somewhat by the presence of the magnetic field. Suppose now that the same total explosion energy in the magnetized and nonmagnetized superbubbles at some epoch is shared among the kinetic, thermal, and

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<tr>
<td>G1.01</td>
<td>Gaussian(^b)</td>
<td>10(^a)</td>
<td>1</td>
<td>7b</td>
</tr>
<tr>
<td>H0.01</td>
<td>~1.00(disk)</td>
<td>~10(^{0})(disk)</td>
<td>0</td>
<td>8, 9, 12, 14, 15</td>
</tr>
<tr>
<td></td>
<td>~0.01(corona)</td>
<td>~10(^{0})(corona)</td>
<td>1</td>
<td>10, 11, 12, 14, 15</td>
</tr>
<tr>
<td>H1.01</td>
<td>~1.00(disk)</td>
<td>~10(^{0})(disk)</td>
<td>2</td>
<td>12</td>
</tr>
<tr>
<td></td>
<td>~0.01(corona)</td>
<td>~10(^{0})(corona)</td>
<td>2</td>
<td>12</td>
</tr>
</tbody>
</table>

\(^a\) \( n = n_0 \exp\left(-z/H_o\right) \) with \( H_o = 100 \) pc and \( n_0 = 1.0 \text{ cm}^{-3} \).

\(^b\) \( n = n_0 \exp\left(-z^2/2H_o^2\right) \) with \( H_o = 100 \) pc and \( n_0 = 1.0 \text{ cm}^{-3} \). \( R = p_{\text{tot}}/p_{\text{mag}} \) is assumed to be constant in all models.
Fig. 4.—Same as Fig. 2, but for a model with magnetic field (model E1). The magnetic field lines are also displayed by solid lines in the right panels. The plasma $\beta = p_{\text{pl}}/p_{\text{mag}}$ is initially uniform ($\beta = 1$) and $B_z \approx 4.5 \mu G$ at $z = 0$. Elapsed times are (a) 5.2 and (b) 10.6 Myr, respectively.
magnetic energies within the same volume enclosed by the forward shocks in each case. If so, then the reduced postshock velocity mentioned above implies that $P_{\text{ion}}/P_1$ must actually be a little higher for the magnetized case. Since $M_{1,1}$ increases with increasing $P_{\text{ion}}/P_1$, this explains the numerical result that the forward shock travels upward faster when $B \neq 0$.

For a given $\beta_1$, the existence of a shock requires that $M_{1,1} > (\gamma + 2/\beta_1)^{1/2}$ [i.e. $v_{\text{sh}} > (c_{s,1}^2 + v_{A,1}^2)^{1/2}$, where $c_{s,1} = \gamma P_1/\rho_1$ is the preshock sound speed and $v_{A,1} = (B_1^2/4\pi \rho_1)^{1/2}$ is the preshock Alfvén velocity.] In other words, the shock must be both supersonic and super-Alfvénic in order to be a shock in this case. Hence, since we have assumed $\beta_1 = \text{constant in } z$ here, we expect that, as the shock travels in the $z$-direction, it will continue to be both supersonic and super-Alfvénic, as long as it is so at the point at which it
begins to accelerate along the z-axis. We note, however, that if, instead, \( \beta_z \) decreases with increasing \( z \), this condition can be violated.

In that case, the MHD shock weakens to become a nonlinear magnetoacoustic wave, but the qualitative description given above still applies.

The broadening and the deceleration of the contact discontinuity observed in our numerical results is explained in terms of the magnetic pressure, or \( J \times B \) force. As the shock wave propagates, the magnetic field is also swept up by the expanding plasma following the shock front, thus forming a “void” of very weak magnetic fields (a diamagnetic bubble) in the center, and a region of higher magnetic field strength in the contact surface, where the strong magnetic pressure and tension tend to pull ejected mass back as shown by Kulsrud et al. (1965). We can see this in Figure 3: \( p_{\text{mag}} \) (dashed line) is comparable to \( p_{\text{gas}} \) (solid line) there, and the density distribution is lower and less steep beyond this line compared with that in nonmagnetized model (Fig. 3).

We note that, according to Figure 6, the shock front is accelerating in the \( z \)-direction for both the nonmagnetized and magnetized cases throughout the simulations. This means that the pressure of the gas outside the shock is negligible, relative to the postshock pressure even after the shock travels several 100 pc, so we expect the shocks blowout in both cases. However, the blowout of thermal energy and mass contact discontinuities are significantly inhibited by the magnetic field.

### 3.3. Effects of Cooling

In general, the radiative cooling rate is a strong function both of the density and the temperature and is thus rather sensitive to the detailed structure of superbubbles. In our essentially two-dimensional configurations, however, it is hard to evaluate the cooling effect, because we cannot follow the expansion of the superbubble in the \( y \)-direction (perpendicular to the calculated \( x \), \( z \)-plane).

We therefore will not put the cooling term in the energy equation, but instead adopt a simple approach to study the effects of cooling, i.e., change the adiabatic index from \( \gamma = 5/3 \) to 1.05. In the case of \( \gamma = 1.05 \), the temperature is nearly constant in a region of compression. For such a nearly isothermal gas, compression implies “cooling,” or energy loss, relative to the same compression of an adiabatic gas. From the simulations for the nonmagnetized case, we find that the high-density shell represents a very efficient “cooling” of the flow, in fact, and so the upward motion is suppressed to some degree. In the magnetic model, however, the difference caused by a change in \( \gamma \) is fairly small. This is because the magnetic pressure prevents the contraction of the shell so that the cooling is much less efficient.

In the absence of a magnetic field, the gas in the dense shell which bounds the hot interior would cool radiatively in a time short compared to the overall flow time and at nearly constant pressure, leading to a further increase of the density in the shell. In the case with a magnetic field, however, this radiative cooling cannot result in such isobaric compression of the gas in the shell, since the magnetic field strength in the shell is already large enough to make magnetic pressure dominate over gas pressure there. In that case, the radiative cooling of the gas in the shell will not affect the dynamics as we have calculated it. Our neglect of radiative cooling is partly justified, therefore, in the magnetized case. There is, of course, the possibility that thermal instability could result in small-scale inhomogeneity along the field lines as a result of such cooling.

Finally, since our calculations are two-dimensional, the adiabatic expansion cooling of the gas obeys \( T \sim V^{1-\gamma} \), with \( V \) increasing only as fast as the increase of the two-dimensional cross-sectional area of the expansion, which is not as fast as the increase of the volume of a three-dimensional explosion. Hence, \( T \) tends to stay higher in our two-dimensional calculations than would be expected in a truly three-dimensional explosion. This, in turn, also makes radiative cooling less important and makes our neglect of radiative cooling inside the bubble self-consistent.

### 3.4. Gaussian Model without and with Magnetic Field: G0 and G1

The evolution of superbubbles in a Gaussian atmosphere is illustrated in Figures 7a and 7b for the cases without magnetic field and with magnetic field, respectively. The most striking feature is a blowout of the bubbles at the time \( \sim 6 \times 10^8 \) yr, which is a result of the rapid decrease of the density in high \( z \) (see Mac Low, McCray, & Norman 1989). We note that only a small fraction of the mass is actually accelerated upward. Moreover, the breakup of the shell follows the Rayleigh-Taylor instability, although coarseness of the grid spacings makes it difficult to resolve this phenomenon in Figure 7. The effects of the horizontal magnetic field in this case are similar to those we already discussed in the exponential models (see § 3.2). Here again, we see the acceleration of the shock front and the broadening and the deceleration of the contact surface, relative to the nonmagnetized case.

### 3.5. Model with Galactic Halo: H0, H1, and H2

As more realistic models, we include finite pressure in the halo regions in models H0, H1 and H2. Results for the density contours, the magnetic field lines, and the velocity vectors are displayed in Figures 8 and 10, and the density and the pressure distribution along the \( z \)-axis are plotted in Figures 9 and 11 for the nonmagnetized (H0) and magnetized (H1: \( \beta = 1 \) everywhere) models, respectively. We see two contact surfaces; one is a result of an explosion (same as seen before), and the other comes from the initial density discontinuity. It might be noted that, although we put a high pressure halo above the disk here, the disturbance travels faster than in the previous model in which the pressure is practically zero in the halo. This is a case in which, even though \( \beta_z = \text{constant with } z \), the high temperature of the halo implies high sound speed \( c_{s1} \), so that the condition of shock existence, \( M_{1,1} > (\gamma + 2/\beta_z)^{1/2} \), is violated once the shock reaches the halo. As such, the shock cannot continue to exist in the halo. The fact of constant \( \beta_z \) means that both the sound speed, \( c_{s1} \), and the Alfvén speed, \( v_{A1} \), are high in the hot halo (i.e., \( \beta_z = c_{s1}^2 / v_{A1}^2 = \text{constant} \)). As a result, the shock weakens upon entering the halo to become a nonlinear magnetoacoustic wave. In Figure 12 we plot the positions of the wave front (WF) and of the contact discontinuities (CD) for models H0, H1, and H2 (where \( \beta = 0.5 \) everywhere initially). According to Figure 12, even though the magnetic field causes the disturbance to propagate faster in the \( z \)-direction than without the field, the contact discontinuity, which bounds most of the mass and thermal energy from the SN shock heating of disk gas, is significantly decelerated in the magnetized case.
Fig. 7a

Fig. 7b

Fig. 7.—Same as Fig. 2 but for the Gaussian model (a) without magnetic field, a model G0 and (b) with magnetic field, model G1, respectively. The magnetic field strength in model G1 is $B_z = 4.5 \mu G$ at $z = 0$ and decreases with $z$ as a Gaussian. Elapsed times are 8.2 Myr for both cases.
Fig. 8.—Same as Fig. 2, but for the model with the halo H0. Elapsed times are (a) 3.0 and (b) 4.8 Myr, respectively.
For comparison, we calculated models with the same density distribution as models H0 and H1, but constant temperature, $T = 10^4$ K, i.e., small pressure in the halo. The results are consistent with the exponential model, E0, and not with this (H0). Thus we conclude that the larger disturbance propagation speed is due to the high sound and Alfvén velocities in the high-temperature halo.

Note that we assumed that the pressure in the halo is roughly independent of $z$ in models H0 and H1. Nevertheless, the outermost disturbance defined by WF in both H0 and H1 appear to be capable of blowing out, since no deceleration is evident out to several hundred parsecs, although the blowout of the contact discontinuity may be prevented by the magnetic field. If we allow the halo pressure to decrease with height, we expect WF to accelerate toward high $z$, and the contact discontinuity can also blowout even with the magnetic field. We note that the acceleration due to gravity is small compared with that due to the pressure gradients, so gravity has only a minor effect.

When we include the horizontal magnetic field, the shape of the shell is much more elongated in the horizontal direction, whereas that of the nonmagnetized model is rather spherical. This suggests that the observed elongated shapes of some supernova remnants might be evidence of strong magnetic fields (discussed later).

4. SUMMARY AND DISCUSSION

4.1. Schematic MHD Superbubble

The qualitative results of our numerical simulations of magnetic superbubbles are summarized in the schematic diagram in Figure 13 (corresponding to model H1). The magnetic fields which were originally located in the halo are now swept up by the outer contact discontinuity (OC) and form a “magnetic wall” just above OC, in which the field strength is several times larger than the initial field strength. The dense shell between the magnetic wall and the hot bubble is a compressed, swept up shell, not a shocked shell. Although gas motion through the magnetic wall is largely suppressed, a small amount of gas passes through the wall and contributes to the propagation of the MHD forward shock (FS) or waves, carrying a certain fraction of the energy with it. This is a difference between a magnetic wall and a real wall (caused by a density discontinuity), which suppresses motions through the boundary surface completely (see, e.g., Mineshige & Shibata 1990). The reverse shock (RS) was not clear in our calculations, because we continually supply mass and energy (see discussion in § 3.1). The basic shock-bubble structure displayed in Figure 13 is also relevant to the exponential models (E) and the Gaussian models (G), although the outer contact discontinuity is missing in these models because of the absence of the initial density discontinuity.
Fig. 10a

Fig. 10b

Fig. 10.—Same as Fig. 4 but for the model with magnetic field, model H1. Elapsed times are (a) 3.0 and (b) 4.8 Myr, respectively.

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4.2. Intermediate MHD Shocks

One of the most intriguing features of our numerical MHD results is the apparent development of a stable "intermediate" MHD shock at the outer edge of the superbubble along its sides, where the angle between the preshock magnetic field direction and the velocity vector of the preshock gas in the frame of the outward moving shock is relatively small. Such shocks are characterized by a reversal of the direction of the component of the field parallel to the shock plane which occurs across the shock. This is well illustrated by the lower right portion of the right panel in Figure 10b, for example.

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FIG. 12

Fig. 12—Same as Fig. 6, but for models H0, H1, and H2

Fig. 13—Schematic view of the evolution of a magnetic superbubble (corresponding to model H1). The symbols "FS" and "RS" stand for the forward and the reverse shock, and "OC" and "IC" represent the outer and the inner contact discontinuities, respectively.

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Contrary to the usual belief that MHD intermediate shocks are unstable or “extraneous” (e.g., Jeffrey & Taniuti 1964; Kantrowitz & Petschek 1966), it has recently been suggested by numerical solutions that at least some intermediate shocks may be stable and admissible (Wu 1987). Intermediate shocks differ from fast and slow MHD shocks in that, unlike the latter shocks, there is not an analogous, compressive, intermediate wave that may steepen to form the shock. However, as demonstrated recently by two-dimensional MHD numerical simulations (Steinolfson & Hundhausen 1988; Shibata et al. 1989b), intermediate shocks may form in environments where adjacent (crossflow) regions have the opportunity to interact with and influence each other. The superbubble simulations presented here offer another example of the conditions which can lead to a persistent intermediate shock.

As mentioned above, a primary difference between intermediate shocks and the other two types is that the component of the magnetic field parallel to the shock plane reverses direction across the former. Whereas the magnetic field strength increases across fast shocks and decreases across slow shocks, it may either increase or decrease across an intermediate shock.

### 4.3. Energy Balance

Figure 14a and 14b display the energetics of models E0 and E1, and H0 and H1, respectively. Since we are only concerned with the transformation of the explosion energy, we subtract the initial thermal energy and the magnetic energy of the background medium in the volume enclosed by the disturbance. Initially, all the explosion energy is in the form of the thermal energy. The total energy increase with time, because a new supernova explosion is assumed to occur every 0.3 Myr. For the exponential cases, E0 and E1, the thermal energy fraction decays to $E_{th}(E0) \sim 60\%$, and $E_{th}(E1) \sim 65\%$ by $t = 10$ Myr, almost linearly with time. The kinetic energy fraction decays during this time from $\sim 20\%$ at $t \sim 1$ Myr to $E_{kin}(E0) \sim 40\%$ and $E_{kin}(E1) \sim 30\%$ by $t = 10$ Myr, also linearly with time, and will eventually escape from the system by the shock. The magnetic energy fraction in case E1 grows only modestly, never exceeding $\sim 5\%$ by 10 Myr. In the disk halo cases, the nonmagnetized case yields similar fraction for $E_{th}$ and $E_{kin}$, although without the linear time dependence. However, in the magnetized case, the magnetic share grows almost linearly in time to $E_{mag} \sim 20\%$ equal to $E_{kin}$. Note the difference between nonmagnetized model H0, in which 40% of energy goes into kinetic energy. This indicates that the presence of magnetic fields causes significant reduction in kinetic energy, and the remaining fraction of input energy is converted and stored in the magnetic fields without going out of galactic disk halo region.

**Fig. 14a**

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To see how mass and energy are carried upward, we illustrate mass flow rates,

$$\dot{M}(z) = y_{\text{exp}} \int \rho v_x(x, z) dx,$$

(upper) and energy flow rates

$$\dot{E}(z) = y_{\text{exp}} \int \left[ \frac{\gamma}{\gamma - 1} \rho v_x^2 + \frac{1}{2} \rho v_x (v_x^2 + v_y^2) + \frac{B_x}{4\pi} (v_x B_x - v_y B_y) \right] dx,$$

(lower) as functions of $z$ for models H0 and H1 in Figures 15a and 15b, respectively. Here again $y_{\text{exp}} = 500$ pc (see § 2.3). We note that upward flow of mass and energy is significantly suppressed in magnetized model (see Fig. 15b), but still a small fraction of mass and energy leaks from the magnetized hot bubble and goes upward.

### 4.4. Three-dimensional Effects

Since the magnetic fields have essentially three-dimensional structure, whereas we only performed the two-dimensional simulations, we should add some remarks on the three-dimensional effects of magnetic fields. If the magnetic field is uniform both in direction and in strength, then the suppression of the motion across the field lines is less efficient because of the interchange instability (or the interchange mode of the magnetic Rayleigh-Taylor instability; for a review, see Hughes & Proctor 1988). This is an instability in which, as a result of the exchange between “light” magnetic field lines and “heavy” elements (note that the direction of the acceleration is upward), the matter is less affected by the magnetic pressure, and hence can go up further. When the interchange instability occurs, the magnetic field will form filament structures.

The interchange mode is stabilized by the presence of magnetic shear (nonuniformity in the direction of the field). In such a case, however, other kinds of instability which have larger growth rates appear: e.g., Parker instability (or the buoyancy instability in gravitating systems), or the ballooning instability (buoyancy instability in the presence of an effective gravity caused by the deceleration of a system or by the curvature effect in the case of the magnetic confinement of a high pressure plasma; for a review,
5. APPLICATIONS

5.1. Galactic Fountain

In the Galactic fountain model of Shapiro & Field (1976), gas in the Galactic disk which is heated by successive supernova remnant shocks escapes as hot gas into the halo where it cools radiatively as it flows, until cooler, denser clouds form by thermal instability. These clouds eventually lose buoyancy and fall ballistically back toward the Galactic plane, possibly explaining the high-velocity clouds seen in 21 cm emission. Support for this picture has subsequently come from the observation of UV absorption lines of highly ionized atoms like C iv, Si iv, and N v due to gas more than 1 kpc above the Galactic plane (cf. Savage 1987) and of UV emission lines of C iv and O iii, such as expected from cooling Galactic fountain gas (Martin & Bowyer 1990). The reader is referred to Spitzer (1990) for a recent review of this subject and for references to other, more recent, theoretical work which has been done on it.

Of particular interest to us here is the effect of a localized sequence of supernova explosions such as would result from the evolution of an OB association. In addition to the possibility that such “clustered” supernovae might emplace the observations of supershells and superbubbles, they are also a potential source of the supernova-remnant shock-heated gas which escapes into the halo in the Galactic fountain model if the overlapping explosions “blow out” of the disk. This has been investigated previously as a purely hydrodynamical (i.e., nonmagnetized) phenomenon in numerical simulations by, for example, Tomisaka & Ikeuchi (1986), Tenorio-Tagle, Bodenheimer, & Różycka (1987), and Mac Low, McCray, & Norman (1989). These simulations demonstrate that superbubbles which are driven by a high enough supernova energy injection rate in an appropriately plane-stratified interstellar medium without a magnetic field do indeed “blow out” of the Galactic disk, thereby venting hot gas and energy into the halo as envisioned in the Galactic fountain model.

The MHD simulations we have presented here, however, suggest that such overlapping explosions which are able to “blow out” of the disk in the absence of a magnetic field can be significantly inhibited by the presence of a horizontal magnetic field in the disk and halo of strength comparable to or greater than the average field strength observed in the galactic disk ($B \approx 5 \mu G$). We find that, in the presence of such a horizontal magnetic field, a superbubble which can only marginally blow out of the disk when $B = 0$ everywhere will have much of the swept-up mass which it would otherwise vent upward into the halo decelerated by the magnetic field. Moreover, although the outer disturbance caused by the sequential supernovae actually propagates into the halo as an MHD shock or nonlinear wave faster in the magnetized case, the halo gas thus heated by this disturbance is generally heated only weakly (i.e., to less than $10^3$ K).

In short, such a magnetized superbubble may be better as an explanation of superbubbles and supershells than as a “vent” or “chimney” through which to channel hot gas into the halo as in the Galactic fountain model. This suggests that, if the Galactic fountain model is correct in assuming that supernova-heated gas in the disk escapes into the halo, then either: (1) the disk magnetic field has a scale height which is smaller than that of the gas pressure, so that $p_f(z) > 1$ or even $\gg 1$ as $z$ increases into the halo, (2) superbubbles which blow out into the halo do so in regions of below average field strength or of vertical field line geometry (note: superbubbles themselves create some vertical field components out of an initially horizontal field); (3) some superbubbles occur which are well in excess of the condition for marginal blow out without a magnetic field, so even a 4.5 $\mu G$ uniform, horizontal field extending to kiloparsec heights above the Galactic plane cannot prevent blow-out; or (4) the shock heating and upward flow of gas in the halo implied by the Galactic fountain model depends upon the occurrence of supernovae significantly above or below the Galactic plane.

5.2. Supershells

The magnetic fields swept up on the shell (near the contact surface) around an expanding hot bubble in our simulations may be relevant to the magnetic fields observed in H i shells (Heiles 1984, 1989). In fact, in the case of $\gamma = 1.05$ (i.e., nearly isothermal in the shell), we find $B \approx 10 \mu G$ and $P_{mag}/P_{gas} \approx 3-10$ in the shell. These values are consistent with observed values: $B \approx 6.4 \mu G$, and $P_{mag}/P_{gas} \gg 1$ (Heiles 1989). A similar view on the relationship between hot bubbles and magnetic fields has been suggested by Hayakawa (1979) to explain the observed anticorrelation between soft X-ray enhancement and radio continuum emission.

5.3. Superbubbles in the Galactic Center and Starburst Galaxies

It has been suggested that there may be a few hot bubbles created by successive supernova explosions originating from star bursts in the center of our galaxy (Umemura et al. 1988; Sofue 1989). According to Tsuboi et al. (1989), vertical magnetic fields observed in the center of the Galaxy may be a result of interaction between the initially horizontal magnetic field and multiple superbubbles. In fact, even one superbubble creates vertical components of magnetic fields out of horizontal fields as shown, e.g., in Figure 7b.

Recent observations have revealed that hot outflows are emanating from the central regions of some spiral galaxies such as NGC 3079 (Hummel et al. 1983; Duric et al. 1983; Irwin et al. 1987) and starburst galaxies such as M82 (Nakai et al. 1987; Sofue 1988). Some of the outflows show loopylike structure (e.g., NGC 3079), which may be due to the interaction of expanding superbubbles and horizontal magnetic fields.

5.4. Application to Supernova Remnants

The nonspherical configuration of bubbles found in our adiabatic simulations may be applied to nonspherical remnants, since there is no intrinsic scale in the adiabatic models. Kesteven & Caswell (1988) found that the dense shells and filaments in G296.5+10.0 and G327.6+14.6 are elongated with a barrel shape (cylindrical symmetry), presumably along magnetic field lines.
Landeker et al. (1982) interpreted the origin of the asymmetric shape of the supernova remnant VRO 42.05.01 as the hot cavity. However, this might also be caused by strong magnetic fields (Mineshige & Shibata 1990). A three-dimensional MHD simulation and detailed comparison with the observations would be of some interest in this regard.

After this paper was initially submitted, independent simulations in two-dimensions by Tomisaka (1992) for a uniform density gas were performed. These simulations include radiative cooling, but confirmed our assumption that cooling is unimportant for interior hot gas in superbubbles.

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REFERENCES

Hayakawa, S. 1979, PASJ, 31, 621
Horiiuchi, T., Matsumoto, R., Hanawa, T., & Shibata, K. 1988, PASJ, 40, 147
Ikeuchi, S. 1988, Fund. Cosmic Phys., 12, 4
Priest, E. R. 1982, Solar Magnetohydrodynamics (Dordrecht: Reidel)
Roache, P. J. 1972, Computational Fluid Dynamics (Albuquerque: Hermosa), chap. 5
Shapiro, P. R., Mineshige, S., & Shibata, K. 1993, in preparation
Shibata, K., & Uchida, Y. 1985, PASJ, 37, 31
---. 1990, PASJ, 42, 39
Sofue, Y., & Fujimoto, M. 1989, preprint
Spitzer, L. 1978, Physical Processes in the Interstellar Medium (New York: John Wiley & Sons)
---. 1990, ARA&A, 28, 71
Steinolfson, R. S., & Hundhausen, A. J. 1988, preprint
Tomisaka, K., & Ioku, S. 1986, PASJ, 38, 697