Starburst-driven Mass Loss from Dwarf Galaxies: Efficiency and Metal Ejection

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ABSTRACT

We model the effects of repeated supernova explosions from starbursts in dwarf galaxies on the interstellar medium of these galaxies, taking into account the gravitational potential of their dominant dark matter haloes. We explore supernova rates from one every 30,000 yr to one every 3 million yr, equivalent to steady mechanical luminosities of $L = 0.1 - 10 \times 10^{38}$ ergs s^{-1}, occurring in dwarf galaxies with gas masses $M_g = 10^6 - 10^9 M_\odot$. We address in detail, both analytically and numerically, the following three questions:

1. When do the supernova ejecta blow out of the disk of the galaxy?
2. When blowout occurs, what fraction of the interstellar gas is blown away, escaping the potential of the galactic halo?
3. What happens to the metals ejected from the massive stars of the starburst? Are they retained or blown away?

We give quantitative results for when blowout will or will not occur in galaxies with $10^6 \leq M_g \leq 10^9 M_\odot$. Surprisingly, we find that the mass ejection efficiency is very low for galaxies with mass $M_g \geq 10^7 M_\odot$. Only galaxies with $M_g \lesssim 10^6 M_\odot$ have their interstellar gas blown away, and then virtually independently of L. On the other hand, metals from the supernova ejecta are accelerated to velocities larger than the escape speed from the galaxy far more easily than the gas. We find that for $L_{38} = 1$, only about 30% of the metals are retained by a $10^9 M_\odot$ galaxy, and virtually none by smaller galaxies. We discuss the implications of our results for the evolution, metallicity and observational properties of dwarf galaxies.

Subject headings: Hydrodynamics – Shock waves – Galaxies: Dwarfs – Numerical Methods

1. Introduction

Hierarchical models of structure formation suggest that dwarf galaxies are the building blocks of larger galaxies, merging at high redshift to form the distribution of galaxies we see today. Dwarfs, and even smaller stellar systems, are also responsible for the origin of some of the Lyman α forest absorption features observed in the spectra of distant quasars (Fransson & Epstein 1982, Wang 1995, Ciardi & Ferrara 1997). It is therefore crucial to understand both their formation and evolution. In turn, it has recently become clear that supernova-driven winds play a crucial role in the evolution of such objects since they regulate the mass, metal enrichment, and energy balance of the interstellar medium (ISM) in these galaxies. These aspects have been investigated in a number of theoretical papers (Larson 1974; Dekel & Silk 1986; Silk *et al.* 1987; Vader 1986, 1987; and Ferrara & Tolstoy 1998, hereafter FT) both in local and cosmological contexts.

The observational evidence in support of the existence of outflows from dwarf galaxies has grown rapidly in recent years. Meurer *et al.* (1992) found an expanding bubble of ionized gas with $\sim$ kpc size and velocity of order 100 km s^{-1} in the active dwarf galaxy NGC1705, and Martin (1996) found evidence for superbubbles in I Zw 18 in which the ionized gas expands with velocities of several tens of km s^{-1}. Similar results have been obtained by Della Ceca *et al.* (1996; see also Heckman *et al.* 1995) from ASCA observations of NGC 1569, a star-forming dwarf galaxy in which the authors have been able to demonstrate the presence of diffuse hot gas ($T \sim 0.7$ keV) very likely heated by supernova explosions. Bomans, Chu, & Hopp (1997) found X-ray emission from hot gas within a supergiant shell in the dwarf irregular NGC 4449, which also has a large-scale radio synchrotron halo (Klein *et al.* 1996). In their search for outflows in dwarf galaxies Marlowe *et al.* (1995) conclude that the outflow phenomenona are relatively frequent in centrally star-forming objects; in addition they point out that these outflows tend to be generally oriented along the galaxy minor axis. On a smaller scale, Puche *et al.* (1992) in their VLA HI study of dwarf galaxies, have detected low-density bubbles surrounded by a shell expanding with velocity $\sim$ 15 km s^{-1}. An absorption line study (Bowen *et al.* 1996) from low column density gas in the dSph galaxy Leo 1 has given upper limits on the amount of hydrogen present in the halo of such galaxy ($M \leq 10^4 M_\odot$). This low value seems to be in contradiction with the expectations from simple blowaway models in which the gas content of the galaxy has been ejected into the halo. However, as the authors stress, the possibility that a significant fraction of *hot* gas exists along the line of sight cannot be ruled out. (Alternatively, even if the galactic gas had indeed been ejected, it could have been stripped either by ram pressure against the intergalactic medium (IGM) or by a galactic encounter.)

In many cases these observed gas velocities clearly exceed the escape velocity from the potential well of the parent galaxy, and therefore this material will be injected into the IGM and lost to the galaxy. However, it is difficult to determine what fraction of the ambient medium is indeed brought to such high speeds. Mass can be ejected from a galactic disk through the "blowout" of a superbubble driven by multiple supernova remnants, carving a hole through the galactic disk (Mac Low & McCray 1988). Gas accelerated above the escape velocity is then "blown away", and escapes the potential of the galaxy entirely (De Young & Heckman 1994). Determining the conditions under which blowout occurs, and under which most of the ISM can be blown away are necessary steps towards the understanding of the evolution of such objects, and also of the possible relationships between early and late type dwarfs.

In particular, for blowout there is yet no clear answer to the question of what fraction of the gas swept up in the shell is accelerated above the escape velocity. The reason for this is that most previous studies in which blowout has been simulated numerically (Tomisaka & Ikeuchi 1986, 1988; Mac Low *et al.* 1989; Tenorio-Tagle, Rozyczka, & Bodenheimer 1990; Silich 1992; Mineshige, Shibata, & Shapiro 1993, Suchkov *et al.* 1994) have not been followed for long enough evolutionary times. Moreover, these models have been almost totally concerned with superbubbles and starbursts in massive galaxies; thus it is not straightforward to extrapolate their results to dwarfs, which have rather different properties, both in terms of their ISM and, perhaps more important, of their dark matter content.

De Young & Gallagher (1990) did model a dwarf galaxy undergoing a starburst, and concluded that about 1/6 of the mass of the swept-up shell can escape the galaxy. However, their study neglected the presence of dark matter in the galaxy, and assumed a thin gaseous disk with a scale height of 150 pc, as well as suffering from insufficient numerical resolution, leaving their conclusions questionable. Ideally, one would also like to find how the fraction of mass that can escape depends on the metallicity of the gas; the cooling provided by metals can be strongly inhibited in the relatively unpolluted gas of dwarfs, whose metallicity is typically more than one order of magnitude below solar (Skillman, Kennicutt, & Hodge 1989).

A related important point concerns the fate of metals ejected by supernovae exploding in the burst. The perceptive idea introduced by Vader (1986) that outflows can be "metal enhanced", although appealing, needs to be quantitatively examined. De Young & Gallagher (1990) did try to investigate the fate of metals in their model. They concluded that most of the metal-rich supernova ejecta are indeed ejected, but that as much as 1/3 of the ejecta end up within the swept-up shell surrounding the hot bubble. However, this appears to have been due to unphysical diffusion of their tracer particles from the hot interior gas of the bubble into the surrounding cold shell, due to the low numerical resolution of their model. If this result were valid, then the concept of metal-enhanced winds would be much less viable, since the metal-to-gas ratio in the wind could not achieve values much higher than in the ambient ISM of the galaxy. Given the apparent problems with the computation, however, it seems worthwhile to reexamine this issue.

In summary, the present study concentrates on dwarf galaxies, and attempts to answer the following questions:

1. What are the conditions for either blowout or blowaway to occur (see our definitions in § 3)?
2. What fraction of gas escapes the galaxy when they do?
3. What is the fate of metals ejected by the central supernovae?

We investigate these problems taking into account the full structure of dwarf galaxies, by making physically meaningful models for the structural properties of dwarf galaxies in order to determine the characteristics of the ambient medium in which the multiple supernova explosions take place. We combine analytic calculations with numerical hydrodynamic simulations using ZEUS-3D[1], a code using the algorithms described by Stone & Norman (1992).

The plan of the paper is as follows. In § 2 we describe our model for the structure of dwarf galaxies, based on recent data on their dark matter and gas content. In § 3 we describe how we perform our numerical computations. Before presenting the results of these experiments in § 5, we derive in § 4 simple analytical conditions for the occurrence of blowout and blowaway as a function of the mass of the galaxy and the mechanical luminosity of the central supernovae. § 6 contains discussion and a brief summary of the paper.

2. Structure of Dwarf Galaxies

For our purposes, we model a dwarf galaxy as a system made of a visible disk with combined gaseous and stellar mass M_g, and a dark matter halo of mass M_h. We assume that the gas has a density distribution given by $\rho_g(\varpi, z, \phi) = \rho_0 f(\varpi, z)$, where $\rho_0 = \rho(0,0)$, ϖ is the galactocentric radius, z is the vertical coordinate, and $f(\varpi, z)$ is a function

[1]Developed by the Laboratory for Computational Astrophysics at the National Center for Supercomputer Applications, and available for community use by registration at the email address lca@ncsa.uiuc.edu

determined, as described below, by imposing hydrostatic equilibrium on the gas in the total gravitational potential of the dwarf galaxy, $\Phi_t(\varpi, z) = \Phi_h(\varpi, z) + \Phi_g(\varpi, z)$, which consists of halo and disk components.

The behavior of the dark-to-visible mass ratio $\phi = M_h/M_g$ in galaxies has been explored in great detail in a key study by Persic *et al.* (1996). These authors find that ϕ is a function of the galactic mass; using their relations one can easily derive the dependence of this ratio on the visible mass of the galaxy:

$$\phi \simeq 34.7 M_{g,7}^{-0.29}. \tag{1}$$

where $M_{g,7} = M_g/10^7\ M_\odot$. From this equation it is clear that the gravitational potential of dwarf galaxies with $M_g \lesssim 10^9 M_\odot$ is dominated by the dark matter halo; therefore we will neglect the potential due to visible mass Φ_g in all the following calculations involving the total potential Φ_t. The general trend of increasing dark matter fraction with decreasing galaxy mass is indeed consistent with other kinematical studies of dwarf spheroidal galaxies (Mateo 1997) for which velocity dispersions of a significant number of stars have been derived. However, we caution that equation (1) was derived by Persic *et al.* (1996) from a sample of larger galaxies with $M_h \gtrsim 10^{10} M_\odot$ or $M_g \gtrsim 10^9 M_\odot$. Lacking more suitable data, we are extrapolating those results to lower masses. To make this assumption clear throughout the paper we will keep ϕ indicated in a general form as long as possible.

The density profiles of the dark matter haloes remains uncertain. Cold dark matter cosmological models predict that halo density profiles are essentially self-similar with a weak mass dependence (Lacey & Cole 1994; Navarro *et al.* 1997); specifically the dark matter density should decrease as $\rho_h^{CDM}(r) \propto r^{-1}$ at small radii, where $r^2 = \varpi^2 + z^2$. The detailed form of $\rho_h^{CDM}(r)$ is found to be (Navarro *et al.* 1997)

$$\rho_h^{CDM}(r) = 1500 \rho_c \frac{r_{200}^3}{r(5r + r_{200})^2}, \tag{2}$$

where ρ_c is the central density and r_{200} is the characteristic radius within which the mean dark matter density is 200 times the present critical density, $\rho_{crit} = 3H_0^2\Omega/8\pi G = 1.88 \times 10^{-29} h^{-2}$ g cm^{-3}, where $H_0 = 100h$ km s^{-1} Mpc^{-1} is the Hubble constant; throughout the paper we will use $\Omega = 1$. Equation (2) holds down to $r_{min} \sim 10^{-2} r_{200}$ kpc, where the resolution of the simulations becomes inadequate.

Recent observational work, however, appears to disagree with this prediction. Salucci & Persic (1997) have demonstrated convincingly that dark matter halos have a core, *i.e.* a central region of almost constant density, whose size increases with galaxy luminosity both in absolute units and as a fraction of the optical radius. The structure of very small galaxies with mass $\sim 10^7 M_\odot$ is even less well understood. In view of the uncertain state of the art, we calculate the halo gravitational potential by assuming that the density distribution of the halo can be approximated by a modified isothermal sphere (Binney & Tremaine 1987), which is general enough to be appropriate for an idealized situation such as the one presented here, and does reproduce the observed central core. It follows that

$$\rho_h(r) = \frac{\rho_c}{1 + (r/r_0)^2}. \tag{3}$$

The halo mass as a function of radius is then

$$M_h(r) = \int_0^r dr' 4\pi r'^2 \rho_h(r') = 4\pi \rho_c r_0^3 \left(x - \arctan x\right), \tag{4}$$

where $x = r/r_0$. Thus, if $x \gg 1$, which we will show later on,

$$M_h = M_h(r_h) \simeq 4\pi \rho_c r_0^2 r_h, \tag{5}$$

where r_h is an appropriately defined halo radius. Following a common *Ansatz* we take

$$r_h \equiv r_{200} = \left(\frac{3\rho_c}{200\rho_{crit}}\right)^{1/2} r_0, \tag{6}$$

as obtained using the definition of r_{200} given above and equation (5). To proceed further we need a relation between the scale radius, r_0, the central dark matter density ρ_c, and the mass of the halo M_h. Burkert (1995) has shown that the total dark matter inside r_0, given by $M_0 = M_h(r_0) = \pi(4-\pi)\rho_c r_0^3$ (which implies $M_0 \simeq 0.21(r_0/r_h)M_h$), is related to r_0 and ρ_c through the relations

$$M_0 = 7.2 \times 10^7 \left(\tfrac{r_0}{\text{kpc}}\right)^{7/3} M_\odot, \tag{7}$$

$$\rho_c = 2.7 \times 10^7 \left(\tfrac{r_0}{\text{kpc}}\right)^{-2/3} M_\odot \text{ kpc}^{-3}. \tag{8}$$

Substituting for M_0 into equations (7) and (8) we get

$$r_0 = 4 \times 10^{-7} \left(\frac{M_h}{M_\odot}\right)^{3/4} \left(\frac{r_h}{1 \text{ kpc}}\right)^{-3/4} \text{kpc}, \tag{9}$$

$$\rho_c = 5 \times 10^{11} \left(\frac{M_h}{M_\odot}\right)^{-1/2} \left(\frac{r_h}{1 \text{ kpc}}\right)^{1/2} M_\odot \text{ kpc}^{-3}. \tag{10}$$

Substituting the value of r_h from equation (6) into these expressions we obtain the dependence of the halo radius on M_h,

$$r_h = 0.016 \left(\frac{M_h}{M_\odot}\right)^{1/3} h^{-2/3}\text{kpc}, \tag{11}$$

and can substitute back into equations (9) and (10), to find

$$r_0 = 5.3 \times 10^{-5} \left(\frac{M_h}{M_\odot}\right)^{1/2} h^{1/2} \text{ kpc}, \tag{12}$$

$$\rho_c = 2 \times 10^{10} \left(\frac{M_h}{M_\odot}\right)^{-1/3} h^{-1/3} M_\odot \text{ kpc}^{-3}. \tag{13}$$

With these assumptions the gravitational potential of the halo is

$$\Phi_h(r) = 4\pi G \rho_c r_0^2 \left(\frac{1}{2}\log\left(1+x^2\right) + \frac{\arctan x}{x}\right). \tag{14}$$

Thus the galaxy has a circular velocity

$$v_c^2(r) = r\frac{\partial \Phi_h}{\partial r} = \frac{4\pi G \rho_c r_0^2}{x}(x - \arctan x); \tag{15}$$

The rotation curve increases rapidly in the inner parts of the galaxy, already being practically flat at $x = 1.5$. The asymptotic value of the rotation velocity at $x \gg 1$ can then be found by using equations (12) and (13) to be

$$v_c = (11.85 \text{ km s}^{-1})(M_{g,7} h \phi)^{1/3}. \tag{16}$$

We tabulate the circular velocities for galaxies of various masses in Table 1, taking $h = 0.65$ and using equation (1) for the mass ratio ϕ.

We then compute the gas density distribution by solving the steady state momentum equation,

$$(\mathbf{v} \cdot \nabla)\mathbf{v} = -\frac{1}{\rho}\nabla P - \nabla \Phi_h, \tag{17}$$

where $P = \rho c_{s,eff}^2$, ρ is the gas density, $c_{s,eff}^2 = c_s^2 + \sigma^2$ is an effective sound speed, c_s is the gas sound speed, and σ its turbulent velocity. To compute r_0, we again assume $h = 0.65$. In the horizontal direction, the gas is supported against the gravitational pull by both thermal pressure and rotation. In analogy with the dark matter component, we suppose that the gas extends out to a disk cut-off radius, ϖ_*. The latter quantity is obtained from a fit to a sample of dwarfs by FT. They find that the radius-HI mass relation is well approximated by the law

$$\varpi_* \simeq \varpi_0 M_{g,7}^{\alpha} = 3 M_{g,7}^{0.338} \quad \text{kpc}, \tag{18}$$

defining the constants ϖ_0 and α. Cutting the disk at an arbitrary radius does introduce a region out of hydrostatic equilibrium; as a result the disk tends to expand at its own effective sound speed during the simulations. As this is rather a low velocity, it only results in a slight rounding of the sharply cut off outer edges of the disks.

We will use equation (18) thoughout the paper, although one should be aware of the presence of some degree of uncertainty. For the sake of clarity, it might be useful to give the explicit expressions for the relations among the three characteristic radii so far introduced in this Section. These are:

$$\frac{r_h}{\varpi_*} \simeq 1.23(\phi h^{-2})^{1/3} \tag{19}$$

$$\frac{\varpi_*}{r_0} \simeq 15.6 M_{g,7}^{-1/6}(\phi h)^{-1/2} \tag{20}$$

$$\frac{r_h}{r_0} \simeq 19.2 M_{g,7}^{-1/6}(\phi h^7)^{-1/6}. \tag{21}$$

The last equation justifies the approximation $x \gg 1$ used in eq. 5.

Outside ϖ_* we postulate the existence of an IGM whose present density, $\rho_{IGM} = \Omega_b \rho_{crit}$ has been estimated by assuming $\Omega_b = 0.02$ and $h = 0.6$, giving $\rho_{IGM} \sim 1.4 \times 10^{-31}$ g cm^{-3}. When $\rho_g(\varpi, z)$ has decreased to ρ_{IGM}, we set it constant for the remainder of the grid. As this gas is not in hydrostatic equilibrium, it tends to accrete onto the galaxy. This efffect is negligible: if the entire grid were filled with gas of this density, it would have a mass of only $\sim 10^5 M_\odot$, and if it accreted at a constant rate over a typical 200 Myr run, it would give an accretion rate of only $5 \times 10^{-4} M_\odot$ yr^{-1}.

With the above assumptions, the solution of equation (17) provides a reasonably self-consistent initial distribution for the gas in which supernova explosions are simulated. The derived distribution is uniform to 1% in the equatorial plane of the galaxy. A cut along the vertical axis through the galactic center shows that a Gaussian distribution matches the actual distribution very well.

However, in order to provide in the simplest possible manner a tractable analytic estimate for the order of magnitude of the scale height H, $n_0 = \rho_0/\mu m_h$, and N_H, the gas column density, we make the assumption that all the mass is concentrated in a thin disk, producing a constant gravitational acceleration $g = 2\pi G \Sigma_t$, where $\Sigma_t = (M_d + M_h)/2\pi\varpi_*^2$ is the total matter surface density, using the prescription for the ratio of dark to visible matter ϕ given by equation (1). Such a potential would generate an exponential gas distribution with parameters given by the following formulae:

$$H = \frac{c_{s,eff}^2}{2\pi G \Sigma_t} = \frac{c_{s,eff}^2 \varpi^2}{G M_g (1+\phi)} \simeq 2 \left(\frac{\varpi_0^2}{\phi}\right) M_{g,7}^{2\alpha-1} c_{10}^2 \quad \text{kpc}, \tag{22}$$

$$n_0 = \frac{M_g}{2\pi H \varpi^2 \mu m_h} = 2 \times 10^{-2} \left(\frac{\phi}{\varpi_0^4}\right) M_{g,7}^{2(1-2\alpha)} c_{10}^{-2} \quad \text{cm}^{-3}, \tag{23}$$

$$N_H = n_0 H = 2 \times 10^{20} M_{g,7}^{(1-2\alpha)} \varpi_0^{-2} \quad \text{cm}^{-2}, \tag{24}$$

where, $c_{10} = c_{s,eff}/10$ km s, ϖ_0 is defined by equation (18) in kpc, and μ is the mean molecular weight. We tabulate in Table 1 these values as a function of M_g for the given values of $c_{s,eff}$. We use only n_0 from these values in the simulations, but they are useful for analytic estimates. We note that in this simple approximation, N_H is independent of both ϕ and $c_{s,eff}$ and, given that $1 - 2\alpha \simeq 1/3$, only weakly dependent on the visible mass of the galaxy.

3. Analytical Considerations and Estimates

Before we present the results of the numerical simulations, it is instructive to derive simple analytical relations that approximately give the final fate of the shocked gas and of the galaxy itself.

We distinguish between two potential results of a central starburst. First, in a "blowout", the central supernova explosions blow a hole through the galactic gas distribution, parallel to the steepest density gradient (usually along the rotation axis), accelerating some fraction of the gas and releasing the energy of subsequent explosions without major effects on the remaining gas. Second, in a "blowaway", all, or nearly all, of the ambient ISM is accelerated above the escape velocity and is lost to the galactic potential well. In the following we derive in brief the conditions under which the two processes can occur, as a function of the mechanical luminosity of the starburst, and of the mass of the galaxy. A more extended discussion can be found in FT.

3.1. Blowout

The blowout condition can be derived by requiring that the blowout velocity v_b exceeds the escape velocity of the galaxy v_e. The blowout velocity can be defined as the velocity at a height $z = 3H$ above the galactic midplane, where H is the exponential scale height. Since the velocity of a shock produced by an explosion in a stratified medium decreases down to a minimum occurring at $z \approx 3H$ before being reaccelerated in case of blowout, the above definition of v_b corresponds to the minimum of such a curve. The blowout condition follows from an analysis of the Kompaneets (1957) solution for an explosion in a stratified medium. The explicit expression for v_b has been obtained by FT:

$$v_b = \frac{e^{3/2}}{3^{2/3}} \left(\frac{125}{154\pi}\right)^{-1/6} \left(\frac{L}{\rho_0}\right)^{1/3} H^{-2/3}. \tag{25}$$

We recall that the gas mass for a vertically exponential, horizontally constant density distribution is

$$M_g = 2\pi\varpi_*^2 \mu m_h n_0 \int_0^\infty dz e^{-z/H} = 2\pi\varpi_*^2 \mu m_h n_0 H. \tag{26}$$

Using equations (22), it follows that

$$v_b = 2.7(4\pi GL)^{1/3}(1+\phi)^{1/3} c_{s,eff}^{-2/3}. \tag{27}$$

In general the effective sound speed $c_{s,eff}$ has a contribution from a turbulent velocity, but we make no attempt to model its value, merely assuming that it is constant.

Eq. 27 can now be cast in the convenient form

$$v_b = 92 L_{38}^{1/3}(1+\phi)^{1/3} c_{10}^{-2/3} \quad \mathrm{km\ s^{-1}}. \tag{28}$$

We note that the value of the blowout velocity does not depend on the radius ϖ_*, but only on the the mechanical luminosity of the explosion and the dark-to-visible mass ratio.

Now we can compute the escape velocity v_e at the disk radius for our comparison. It is

$$v_e^2(\varpi_*) = 2|\Phi_h(\varpi_*)| \sim 8\pi G \rho_c r_0^2 \left[\frac{1}{2}\log(1+x_*^2) + \frac{\arctan x_*}{x_*}\right] \quad x_* = \varpi_*/r0. \tag{29}$$

Taking the appropriate value for x_*, as derived from eq. 19, the terms in parenthesis give a factor $\simeq 1.65$ for all galactic gas masses, since it is $x_* \gg 1$. It follows that

$$v_e(\varpi_*) \sim (13.2\pi G\rho_c)^{1/2} r_0 \sim 20 M_{g,7}^{1/3}(\phi h)^{1/3} \quad \mathrm{km\ s^{-1}}, \tag{30}$$

where we've used equation (1) for the relation between the dark-to-visible mass ratio and the galactic mass, and h is the scaled Hubble constant. Again, we remark that v_e is basically independent of ϖ_* as long as $x_* \gg 1$, which implies an almost flat $v_e(\varpi)$. Finally, we can compare v_e to the blowout velocity v_b to find the condition for the blowout to occur:

$$L_{38} > 1.2 \times 10^{-2} M_{g,7} c_{10}^2 h. \tag{31}$$

3.2. Blowaway

We now derive the necessary condition for a starburst to blow away (completely unbind) the ambient gas of the galaxy. We have seen that blowout takes place when the shell velocity at $z = 3H$ exceeds the escape velocity. Following the blowout the pressure inside the cavity drops suddenly due to the inward propagation of a rarefaction wave. The lateral walls of the shell, moving along the major axis of the galaxy, will continue to be accelerated by the interior pressure until they are reached by the rarefaction wave at $\varpi = \varpi_c$, corresponding to a time t_c elapsed from the blowout. After that moment, the shell enters the momentum-conserving phase, since the driving pressure has been dissipated by the blowout. The requirement for the blowaway to take place is then that the momentum of the shell (of mass M_c at ϖ_c) is larger than the momentum necessary to accelerate the gas outside ϖ_c (of mass M_o)at a velocity larger than the escape velocity:

$$M_c v_{\varpi_c} \geq M_o v_e, \tag{32}$$

With some simple assumptions and algebra outlined in FT, one can write the condition for the blowaway to occur:

$$\frac{v_b}{v_e} \geq (\epsilon - a)^2 a^{-2} e^{3/2}, \tag{33}$$

where the radius of the shell in the plane when the pressure is released is $\varpi_b = aH \sim (2/3)H$, and $\epsilon = \varpi_*/H$ is the ratio of the major to the minor axis of the galaxy. Using the expressions given in § 2 for ϖ_* and H, we find

$$\epsilon = 0.43 \frac{\phi}{\varpi_0} M_{g,7}^{1-\alpha} c_{10}^{-2}. \tag{34}$$

Flatter galaxies (larger ϵ values) preferentially undergo blowout, whereas rounder ones are more likely to be blown away. The condition for the blow away to occur, obtained expanding eq. 33 assuming that $\epsilon \gg a$, is then

$$L_{38} > 8 \times 10^{-2} M_{g,7}^{7-6\alpha} \left(\frac{\phi}{\varpi_0}\right)^6 c_{10}^{-10} h. \tag{35}$$

We point out the strong dependence of blowaway on the amount of turbulence present in the galactic ISM: as the value of c_{10} is increased, less powerful explosions can lead to complete disruption of the galaxy. This is because virialized objects with higher effective sound speed tend to be rounder, and hence more fragile with respect to blowaway. The same type of situation might occur if for some reason the distribution of the gas is stretched and becomes less concentrated, as, for example, after the first episode of star formation, or following a galactic encounter. Thus after a first stage of the evolution in which blowout dominates, blow away might occur if the gas distribution can be puffed up enough.

3.3. Final Fate of the Galaxy

From the previous results we can predict the final fate of the galactic gas after a starburst episode. In Figure 1 we show the different regions of the $M_g - L_{38}$ plane in which blowout or blowaway occurs. In this Figure we have used equation (18) for ϖ_*, equation (1) for ϕ, and equation (34) for ϵ. Galaxies with gas masses below $M_{g,7} = 0.04 - 0.1$ undergo blowaway almost independently of the values of the mechanical luminosity of the starburst in the range $L_{38} = 0.1 - 10$. The blowout has a somewhat steeper dependence on $M_{g,7}$: for the luminosities $L_{38} = 0.1 - 10$, blowout can take place up to masses of $\sim 10^8 - 10^{10} M_\odot$, respectively. This, as we will see in the next Sections, is in excellent agreement with the results of our numerical simulations.

4. Initial Conditions of the Simulations

For our numerical models, we use ZEUS-3D, a second-order, Eulerian, astrophysical, gas dynamics code (Stone & Norman 1992) using Van Leer (1977) monotonic advection.

We have implemented equilibrium radiative cooling using the cooling curve given by MacDonald & Bailey (1981), which is a Raymond & Smith (1977) cooling curve extended to lower temperatures. We use an implicit energy equation, implemented with a Newton-Raphson root finder, supplemented by a binary search algorithm for occasional zones where the tabular nature of the cooling curve prevents the Newton-Raphson algorithm from converging. We also include an empirical heating function tuned to balance the cooling in the background atmosphere, but linearly proportional to density, so that it is overwhelmed by cooling in compressed gas (Mac Low *et al.* 1989). This is to prevent the background atmosphere from spontaneously cooling, and may be thought of physically as a crude model for the stellar energy input into the background atmosphere.

We use a tracer field advected with exactly the same algorithm as the density to follow the metal-rich gas ejected by the central energy source. The absence of this field then serves to specify the regions in which cooling should be effective. This is particularly useful in this problem for preventing spurious cooling of the bubble interior due to density numerically diffused off the thin, dense shell of swept up gas surrounding the bubble. In order to maintain a sharp interface at the edge of the hot bubble, we use the method suggested by Yabe & Xiao (1993), which consists of advecting a function $f(c)$ of the tracer field rather than the tracer field c itself. We follow Yabe & Xiao in using the function

$$f(c) = \tan[0.99\pi(c - 0.5)], \qquad c = 0.5 + [\arctan f(c)]/(0.99\pi). \tag{36}$$

With this definition, c will show a sharp transition at the interface if it has initial values of 0 and 1, even if $f(c)$ becomes quite smooth due to numerical diffusion. The factor of 0.99 is used to avoid infinite values for $c = 0$ or $c = 1$. This formulation of the tracer field strongly reduces numerical diffusion of ejected material into the cold shell as compared to the tracer particles used by De Young & Gallagher (1990).

The gas distribution of the galaxy is set up in hydrostatic equilibrium with the dark halo potential given in equation (14, where ρ_c is given by equation (7). To find r_0, we use equation (12), taking $h = 0.65$, and proceeding as described following those equations.

The central energy source is set up as a constant luminosity wind driven by a thermal energy source. We note that Strickland & Stevens (1998) have shown that a single central energy source will more efficiently transfer energy to the ambient gas than multiple smaller clusters with the same total luminosity, as we describe in more detail in § 6. A source region of radius five zones and volume V is set up initially in pressure equilibrium with the background atmosphere, but at a temperature a factor $\eta = 1000$ larger than that of the background atmosphere, and density correspondingly lower by a factor of $1/\eta$. Every timestep we then added energy to the source region at a rate $\dot{E} = L/V$, and mass at a rate $\dot{M} = \rho_0 \dot{E}/\eta e_0$, where e_0 is the midplane energy density of the background atmosphere, maintaining roughly constant specific energy in the source region, and generating a supersonically expanding wind. The amount of mass in this wind is significantly greater than that expected from supernova ejecta alone, as we are trying to also account for mass that would have evaporated off the shell walls had we included thermal conduction in our model. This is only a very rough approximation; as we discuss below, the resulting temperatures are, in fact, somewhat cooler than they would be realistically, which gives us a good lower bound on the metal ejection rates that we compute. We chose this energy input phase to last 50 Myr in our models, the typical lifetime of the least massive star able to become a supernova (e.g., McCray & Kafatos 1987), effectively assuming simultaneous formation of all the stars in the starburst.

We assume azimuthal symmetry and so use in each of our models a two-dimensional grid with the center of the galaxy at the origin, with reflecting boundary conditions along the symmetry axis and along the galaxy mid-plane, and outflow boundary conditions on the other two axes.

5. Numerical Results

We have performed a parameter study at a resolution of 25 pc per zone, varying the gas mass M_g of the dwarf galaxies from 10^6 $M_\odot$ to 10^9 $M_\odot$ (see Table 1 for equivalent rotational velocities and halo masses), and varying the mechanical luminosity of winds and supernovae from the central starbusts from $L = 10^{37}$ ergs s^{-1} to 10^{39} erg s^{-1}. These mechanical luminosities are equivalent to a supernova with energy of 10^{51} erg exploding every 3 million yr to every 30,000 yr, respectively. In this Section we report the results of these models, describing their overall development, computing the efficiency of mass and metal ejection, and making some rough approximations to their observable emission. The results in this section on the occurrence of blowout and blowaway agree very well with the analytic results described in § 3 and summarized in Figure 1.

5.1. Gas Distribution

We begin by showing in Figure 2 a time history of the density distribution of all the models in the parameter study. The first part of the Figure shows the state of the gas at a time of 50 Myr, at the end of the energy input phase. In the higher luminosity models, the internal termination shock can be clearly seen, surrounded by a hot, pressurized region of shocked wind and supernova ejecta. This is in turn surrounded by a colder, radiatively cooled shell of swept-up ambient gas from the galaxy and surrounding IGM. (Another minor effect that can be observed is that the ambient gas at the edge of each galaxy, which is not pressure confined, also begins to expand into the surrounding medium at its sound speed of 10 km/s, producing the rounded ends of the disks.) The larger bubbles have begun to accelerate, and show strong Rayleigh-Taylor instabilities.

Our assumption of azimuthal symmetry has the consequences described in Mac Low *et al.* (1989) for the Rayleigh-Taylor instabilities. In three dimensions, these instabilities produce a characteristic spike and bubble morphology. In azimuthal symmetry, only the central spike can grow, while the other modes are forced into rings. As a result, the central spike is more pronounced than it would be in three dimensions, and the other spikes are less numerous. However, Mac Low *et al.* (1989) showed that changing to an assumption of slab symmetry (*e. g.* using Cartesian coordinates) did not prevent the appearance of a central spike. Although the instabilities do not fully develop because of these effects, they do function to effectively allow hot gas to accelerate beyond the colder, denser, swept-up shells.

After energy input ceases, the holes in the planes of the galaxies recollapse under the influence of gravity and the pressure of the disk gas, except in the extreme cases of low mass and high mechanical luminosity, where the disk gas escapes the potential of the dark matter halo, and is swept completely off the grid at late times. At the same time, the hot bubbles rising above the disks expand out into the surrounding IGM, leaving most of the swept-up shell material behind. After 200 Myr, at the last time shown, the galactic disks have almost completely recovered unless blow-away occurred. Meanwhile, low density gas originating in the central winds and supernovae spreads over regions of tens of kpc. In Figure 3 we show the distribution of metal-enriched wind and supernova ejecta material after 200 Myr for the more massive models. In all cases in which blow-out occurs, this material escapes from the disks of the galaxies; much of it is travelling at high enough velocity to escape from the halo potential as well.

This result should be fairly independent of the details of how the hot gas is traced, so long as hot gas is not allowed to unphysically cool or diffuse into the cold shell. Cooling can happen by numerical diffusion transferring too much mass into the hot interior from the cold shell, as occurred in the models of Tomisaka & Ikeuchi (1986), for example, while diffusion can happen if tracer particles diverge away from the gas they should be following, as occurred in De Young & Gallagher (1990). Physically, the hot gas will not radiatively cool within a dynamical time, and its sound speed is significantly higher than the escape velocity from any of our galaxies, so most of it should indeed escape as is seen in our models.

5.2. Mass and Metal Ejection Efficiencies

We can directly compute the efficiency of mass ejection and retention of metal-enriched gas from our models. The question we ask is how much of each is travelling at speeds higher than the local escape velocity due to the potential of the dark matter halo. In principle, the efficiency of mass ejection is

$$\xi = (M_{\rm esc} + M_{\rm lost})/(M_{\rm i} + M_{\rm sn}), \tag{37}$$

where $M_{\rm esc}$ is the mass on the grid moving at higher than the escape velocity, $M_{\rm lost}$ is the mass lost off the edge of the grid, $M_{\rm i}$ is the mass initially on the grid at $t = 0$, and $M_{\rm sn}$ is the mass injected at the center of the grid by winds and supernovae during the first 50 Myr. To get a good value for ξ we must measure it after most of the gas has been accelerated to its final velocity but before it has left the grid. We chose a value of 100 Myr as a compromise. Although the shells from the lowest mass galaxies have already left the grid, they were moving at above the escape velocity as they went off, so it is reasonable to assume that all mass lost from the grid escaped. We also need to make special provision for the tendency of the relatively large disks in the most massive models to expand sideways off the grid at the sound speed. As no high-speed gas has left the grid in these models, we do this by neglecting $M_{\rm lost}$ for these models. In Table 2 we show the ejection efficiency ξ for each of our models. Only in our most extreme models, with masses of 10^6 M$_\odot$, is most of the mass ejected. In more massive objects, less than 7% of the mass is ejected, usually far less.

The tracer field allows us to separately determine the fate of the metal-enriched gas ejected by the massive stellar winds and SNe of the starburst. We compute the efficiency of ejection of this material ξ_Z by applying equation (37) only to the gas carrying the tracer field. In Table 3 we give ξ_Z for each model. In the more massive galaxies and at lower supernova rates, a significant fraction of the metal enriched gas is retained in the gravitational well of the dark halo, and will eventually fall back on to the galaxy, while at lower masses and higher luminosities, virtually all of the metals escape the grasp of the halo and travel freely into the surrounding intergalactic space.

When estimating the metal ejection efficiency ξ_Z, we have assumed that the metals ejected by massive stars and supernovae are well mixed within the hot interior of the bubble, and are therefore ejected uniformly with this gas. Two questions can be raised here. First, will metals ejected by later supernovae mix with the hot gas of the already existing bubble? A simple order of magnitude calculation (Tenorio-Tagle 1996) shows that the metal diffusion time in the hot bubble gas (mostly consisting of evaporated shell material) is $t_d \sim 10^7 n_{-2} R_{100}^2 T_6^{5/2}$ yr, where $n_{-2} = n/10^{-2}\mathrm{cm}^{-3}$ is the electron density, $R_{100} = R/100$ pc is the radius of the bubble, and $T_6 = T/10^6$ K is the temperature of the hot gas. This time is comparable with, but shorter than the evolutionary timescale of the superbubbles studied in the previous Section.

Second, how will the mixing proceed if SN ejecta are strongly clumped? This might well be the case if Rayleigh-Taylor instabilities act during expansion of the shock through the outer layers of the star (Fryxell *et al.* 1991). Such high-density, high-metallicity clumps could travel almost unimpeded inside the hot, rarefied interior of the bubble, but would be rapidly thermalized when they hit the external shell. In this case the metals would be mixed into the shell, rather than into the escaping hot gas. The conditions under which the clumps can reach the shell, rather than being destroyed in the superbubble interior, have been derived by Franco *et al.* (1993). These authors found that for reasonable density contrasts in the clumps, only the clumps produced by the first 5–10 SNs in the starburst can reach the shell and pollute it with their metal content. The values of $L_{38} = 0.1$–10 and the duration of the energy input phase considered in this paper of 50 Myr imply $\sim$ 15–150 SNs. Thus, ejecta clumping can play a role in the low-luminosity starburts, where we expect that a significant fraction of the metals contained in ejecta clumps ends up not in the hot gas, but rather mixed with the cool shell gas. For these objects our ξ_Z might represent an overestimate; this error should become fairly small for higher luminosity starbursts. It has to be kept in mind, though, that the clumping factor of the ejecta is still far from well-established.

We have checked that our results are reasonably well converged numerically by rerunning one of our models with $M_g = 10^9 M_\odot$ and $L = 10^{38}$ erg s^{-1} at a resolution of 40 zones per scale height, twice that of our standard models. We find that, although details of the shell structure differ slightly, the overall morphology remains almost constant, and the computed efficiencies shift by no more than 10-20%, little enough for our conclusions to remain valid.

5.3. Observables

Finally, we discuss some predictions of observable properties that we can draw from our models. As our assumptions about the microphysics of the gas are quite crude, we are only in a position to make qualitative predictions; however we believe these can nevertheless be useful. Furthermore, we find ourselves in basic agreement with the more careful work performed by Suchkov et al. (1994) for massive starbursts, adding credibility to our approximations.

We begin by computing the temperatures on our grid from the specific energies computed by ZEUS. To do this, we need to assume a composition for the gas, so we take it to be singly ionized with one atom of He for every 10 protons. Figure 4 shows the resulting temperatures (for the more massive models). The hot gas on the grid appears somewhat cooler than physically reasonable. This is most likely due to our replacement of the effects of conduction from the cold shell with the injection of an arbitrary amount of gas at the center of the superbubble as discussed in § 3.

These temperatures can then be used to make rough approximations to the emissivity in X-rays and in optical lines such as Hα. For optical emissivity, we take

$$I_{opt} \propto \rho^2 T^{-1/2}, \tag{38}$$

while for X-ray emission, we take

$$I_x \propto \begin{cases} \rho^2 \sqrt{T} & \text{if } T > T_{min} \\ 0 & \text{otherwise} \end{cases} \tag{39}$$

We chose T_{min} to have an unrealistically low value of 10^5 K in order to more clearly delineate the hot gas emissivity. This choice may be justified for the reasons outlined above, but it should be remembered that these are only qualitative images of the distribution of emissivity.

In Figure 5 we show the optical emissivity, while in Figure 6 we show the X-ray emissivity. Note that these are two-dimensional slices across the galaxy, so they would have to be rotated around the axis and tilted to the line of sight before being directly compared to observations. However, they do clearly indicate the regions expected to be bright in line emission and in X-ray emission. Probably the most important thing to note in this Figure is the clear separation between the relatively dense filaments and the bulk of the hot gas. Deep imaging of optical lines often shows filamentary or bubble-like structures in these dwarf galaxies (e.g. Marlowe et al. 1995; Della Ceca *et al.* 1996); as can be seen from Figure 5 these usually do not trace the outer edges of bubbles but often represent instead fragments of cold shell left behind during the blowout of hot gas to much greater distances.

6. Conclusions

We have explored the effects of stellar winds and SN explosions from starbursts with mechanical luminosities ranging from 10^{37} to 10^{39} ergs s^{-1} on the ISM of dwarf galaxies with mass 10^6 to $10^9 M_\odot$. Specifically, we have addressed in detail the following three questions: What are the conditions for blowout and blowaway to occur, as defined in § 3? What fraction of gas escapes the galaxy in a blowout episode? What is the fate of metals ejected by the massive stars during their lives and in their terminal SN explosions?

We have studied the problem by means of both analytical and numerical techniques, checking — where possible — the agreement between the two types of results. We have found that dwarf galaxies in the above mass range undergo blowout for moderate-to-high luminosity values ($L_{38} \sim 1 - 10$) whereas blowout is inhibited for galaxies with $M \gtrsim 10^8 M_\odot$ if $L_{38} \simeq 0.1$, as shown in Figure 2 and described in Tables 2 and 3. However, the fraction, ξ, of the gas mass of the galaxy lost in such an outflow is surprisingly low in all studied cases: we find that more than a few percent of the mass is lost only in objects with masses as low as $10^6 M_\odot$. In larger objects, the fractions drop to as low as 10^{-6} (Table 2). In the lowest mass objects, the blowaway occurs virtually independently of L. Thus it appears that explosive mass loss is a process involving a threshold determined by the geometry of the gas distribution as well as the depth of the potential well. The coupling between the energy input and the disk gas becomes better for the smaller galaxies as they are more spherical and they need less energy input to begin with for blowaway to occur. It is worth stressing that our results are dependant on the total mass of the galaxy, or the gravitational potential in which the gas is embedded. Although we have tried to model the dark halo carefully, using the best available observational and theoretical results, the precise value of the dark-to-visible mass ratio and the details of dark matter distribution remain quite uncertain. Our general results should be rather solid, but the specific values we quote for efficiency and the mass threshold for blowaway can only be taken as representative.

On the other hand, metal-enriched material from stellar winds and SN ejecta is far more easily accelerated than the ambient gas to velocities above the escape speed. We find that for $L_{38} = 1$ less than 3% of the metal-enriched material is ejected from a galaxy with $M_g = 10^9 M_\odot$ of visible mass, but this fraction increases to 60% for $M_g = 10^8 M_\odot$, and to unity for $M_g = 10^7 M_\odot$ (see Table 3). That is, the smaller galaxies inject virtually all metals produced by the massive stars of the starburst into the intracluster medium or IGM. Stronger starbursts with higher mechanical luminosities eject metal-enriched material even more easily; only for lower starburst luminosities do galaxies retain a large fraction of the heavy elements produced by massive stars.

Our assumption that all energy injection occurs in the central hundred parsecs of the galaxy can questioned, as real starburst galaxies have multiple clusters scattered across the disk, which one might think could be more effective at blowing away the ISM than a single, central energy source. However, Strickland & Stevens (1998) have shown with a simple but elegant argument that our assumption is actually the one leading to the most effective blowaway for a given total luminosity. The gist of their argument can be seen by considering a superbubble in a homogenous medium (Weaver et al. 1977), whose radius $R \propto L^{1/5} t^{3/5}$, but whose kinetic energy $E \propto Lt$. The coupling of a superbubble to the background ISM can be characterized by the energy per unit volume supplied by the superbubble to the ISM $E/V \propto L^{2/5} t^{-4/5}$, which drops with decreasing luminosity. Thus, several small bubbles will be less effective at transferring energy to the ambient gas than one large one, and so multiple clusters will be less effective at blowing away the ISM than our assumed single central cluster.

These results have several implications for the evolution of dwarfs. First, outflows from dwarf galaxies should be strongly metal enriched. Tables 2 and 3 show that the typical metal fraction of the outflow greatly exceeds that of a homogeneously mixed galactic disk. This enrichment occurs because the metals produced by the massive stars in the starburst remain within the hot, shocked cavity gas, which has a high sound speed and so can easily escape the galaxy. On the other hand, it is much more difficult to accelerate the denser, cooler shell of swept-up ambient gas, so that gas remains bound. De Young & Gallagher (1990) found that about 2/3 of the SN ejecta in the chimney escape from the galaxy for a case in which $L_{38} = 0.3$ and $M_g \sim 10^9 M_\odot$. For similar parameters, Table 3 shows that the blowout is prevented, due to the presence of the dark matter halo and a realistic hydrostatic distribution of the gas resulting in a larger scale height (about three times larger than theirs). This suggests that dark matter is key to the evolution of dwarf galaxies and any future models must include it properly.

Our simulations show that after blowout has occurred, a substantial fraction of the hot, metal-rich gas escapes from the galaxy. Metals are then dispersed into the IGM (subject to the caveats discussed in the previous section). As a consequence, it appears that dwarfs could be the major pollutors of the IGM, and certainly have major effects on the environment in which they live. We note the result by Renzini (1997), who finds that most of the iron in clusters appears to reside in the intracluster gas rather than inside galaxies, a clear sign that outflow phenomena are at work in that environment. It is out of the scope of this paper to review the vast literature concerning these issues; for a detailed discussion see FT. We only mention that several authors have already argued that dwarf galaxies might be the most important objects for the IGM enrichment: our calculations provide a solid basis for these previous works and should allow the development of more detailed scenarios.

Starbursts in dwarf galaxies may be traced and studied in detail by X-ray and optical emission-line studies. The results presented here show that huge gaseous halos, with sizes of dozens of kpc are produced, with regions of high x-ray emissivity close to the galactic disk. Relatively cool, dense filaments also occur near the galaxy, well within the external shock, due to shell fragmentation. High spatial resolution spectra of the filaments may be useful to investigate the radiation field in the halo of dwarfs and the escape fraction of ionizing photons from massive stars in the disk, since this cool gas should be predominantly photoionized and hence show up in optical emission lines. Since these filaments are surrounded by the hot gas, the X-ray emission may actually be strongest in regions close to the filaments, where evaporation takes place. Thus, the bulk of the observed X-ray emission may come from these conductive interfaces, in which the gas is far out of ionization equilibrium, and hence emitting strongly. (These effects are not included in Figure 6). The filling factor and the nonequilibrium state of such regions have to be taken into account properly in order to avoid overestimates of the mass and energy involved in the X-ray emission.

The temperature of the hot gas ($T \sim 0.7$ keV) detected by Della Ceca *et al.* (1996) in their observations of NGC 1569 is in qualitative agreement with the ones shown in Figure 4. We have not tried to apply our models to a specific object; an interpretation of the soft X-ray emission from the dwarf galaxy NGC 5253 in terms of a superbubble model has been attempted by Martin & Kennicutt (1995). They concluded that a simple superbubble model with $L_{38} = 1$, exploding in a smooth ambient medium would underproduce the observed X-ray luminosity, $L_x \sim 6.5 \times 10^{38}$ ergs s^{-1}, by a factor ~ 16. As an alternative, they suggest that a clumpy ambient medium generating extra mass of hot gas via cloud evaporation, may help reconcile the above discrepancy. We emphasize that the required hot material could come from the evaporation of the shell fragments seen in our simulations, without the need for a multi-phase medium, whose existence in dwarf galaxies has still to be proven. Due to their large number, dwarf starbursting galaxies may significantly contribute to the X-ray background (Persic *et al.* 1989), although a precise estimate crucially depends on their X-ray luminosity function.

The halos created by the starburst energy injection might also be detectable via absorption line studies towards background objects such as QSOs. Of course this type of experiment is limited by the low column densities in the large, rarefied halo, as well as the availability of suitably placed background objects, which could be a rare occurence for the small dwarf galaxies. A particularly suitable case has been investigated by Bowen *et al.* (1996). The dSph Leo 1 has three QSOs in the line of sight though the halo; however, the authors were only able to put upper limits on the presence of ionized species like CIV, SiV and MgII, but not to exclude the presence of a hotter diffuse component. Similar studies would be particularly important for dIrr and BCD galaxies which are much more active and gas rich and, in general, for normal galaxies.

Finally, we recall that the blowout and blowaway episodes discussed here might be able to carry magnetic field lines along into the halo. It is reasonable to expect that, if a sufficiently strong magnetic field does exist in dwarf galaxies, extended radio-halos should be found around such objects. Radio continuum observations of dwarf galaxies are still scarce; probably the most complete survey has been presented in the pioneering work by Klein *et al.* 1984, who concluded that, on average, blue compact dwarf galaxies exhibit about 10 times higher radio-to-optical luminosity than normal spirals and they have a flatter radio spectrum indicating a weak synchrotron radiation (in turn possibly due to a weak magnetic field). These findings have been largely confirmed by subsequent works devoted to dIrrs (Klein *et al.* 1986) and larger samples (Klein *et al.* 1991). This idea will also apply to larger starburst galaxies.

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Tables

M_g ($M_\odot$)	10^6	10^7	10^8	10^9	10^6	10^7	10^8	10^9
M_h[a]($M_\odot$)	6.8 (7)	3.5 (8)	1.8 (9)	9.1 (10)				
v_c (km s^{-1})	19	34	58	99				
$c_{s,eff}$ (10 km s^{-1})	1.0	1.0	1.0	1.0	0.6	0.6	0.6	0.6
H (kpc)	0.64	0.59	0.55	0.51	0.23	0.21	0.20	0.18
n_0[a](cm^{-3})	5.3 (-3)	1.2 (-2)	2.7 (-2)	6.2 (-2)	1.5 (-2)	3.3 (-2)	7.6 (-2)	0.173
N_H (10^{19} cm^{-2})	1.0	2.2	4.6	9.8	1.0	2.2	4.6	9.8

Table 1: Galactic Parameters

[a]Factors of powers of 10 indicated in parentheses

Visible Mass	Luminosity (10^{38} erg s^{-1})		
($M_g/M_\odot$)	0.1	1.0	10
10^6	0.18	1.0	1.0
10^7	3.5 (-3)	8.4 (-3)	4.8 (-2)
10^8	1.1 (-4)	3.4 (-4)	1.3 (-3)
10^9	0.0	7.6 (-6)	1.9 (-5)

Table 2: Mass Ejection Efficiency, ξ

Visible Mass	Luminosity (10^{38} erg s^{-1})		
($M_g/M_\odot$)	0.1	1.0	10
10^6	1.0	0.99	1.0
10^7	1.0	1.0	1.0
10^8	0.80	1.0	1.0
10^9	0.0	0.69	0.97

Table 3: Metal Ejection Efficiency, ξ_Z

Fig. 1.— Regions of the gas mass (M_g) – mechanical luminosity of the starburst (L) plane in which blowout or blowaway can occur for a values of $c_{s,eff} = 10$ km s^{-1} and $h = 0.65$. The ratio ϵ of the galactic radius ϖ_* to the scale height H is also shown.

Fig. 2.— Density distributions for models with the given initial visible masses M_g and luminosities L at times of (a) 50, (b) 75, (c) 100, and (d) 200 Myr, with the values of density given at the top of each figure. Note that energy input ends at 50 Myr.

Fig. 3.— Distribution of metal-enriched stellar outflow material and supernova ejecta from the starburst energy source at a time of 200 Myr is shown in white, superimposed on the same density distribution shown in Figure 2(d). The material was followed with a tracer field, as described in the text. Note that, in most cases, no enriched gas remains in the disks of the galaxies.

Fig. 4.— A direct computation of the temperatures on our grid at a time of 100 Myr, assuming the gas is ionized with a number fraction of 0.1 of He. These temperatures are qualitatively correct–hot gas in our models should correspond to hot gas physically–but by no means quantitatively correct, as we have neglected important microphysics and injected an arbitrary amount of mass at the center to maintain the superbubble interior density at roughly correct values.

Fig. 5.— A qualitative image of the regions in our models likely to be brighter in Hα and other optical lines, at a time of 100 Myr. We show the emissivity $I_{opt} \propto \rho^2 T^{-1/2}$ on two-dimensional cuts through the galaxies corresponding to our computational grid: these images have not been rotated and projected. Note the prominence of filaments relatively close to the plane of the disk even for extreme blowouts.

Fig. 6.— A qualitative image of the regions in our models likely to be brighter in X-ray emission at a time of 100 Myr. We show an emissivity $I_{opt} \propto \rho^2 T^{1/2}$ for $T > T_{min}$, as discussed in the text. Again, as in Figure 5, these are two-dimensional cuts that have not been rotated and projected. In this image we neglect conductive evaporation off the denser filaments, which may in fact dominate the observed X-ray emission.