Superbubble breakout and galactic winds from disc galaxies

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Accepted 2013 July 10. Received 2013 July 10; in original form 2013 March 11

ABSTRACT
We study the conditions for disc galaxies to produce superbubbles that can break out of the disc and produce a galactic wind. We argue that the threshold surface density of supernovae rate for seeding a wind depends on the ability of superbubble energetics to compensate for radiative cooling. We first adapt Kompaneets formalism for expanding bubbles in a stratified medium to the case of continuous energy injection and include the effects of radiative cooling in the shell. With the help of hydrodynamic simulations, we then study the evolution of superbubbles evolving in stratified discs with typical disc parameters. We identify two crucial energy injection rates that differ in their effects, the corresponding breakout ranging from being gentle to a vigorous one. (a) Superbubbles that break out of the disc with a Mach number of the order of 2–3 correspond to an energy injection rate of the order of $10^{-4}$ erg cm$^{-2}$ s$^{-1}$, which is relevant for disc galaxies with synchrotron emitting gas in the extra-planar regions. (b) A larger energy injection threshold, of the order of $10^{-3}$ erg cm$^{-2}$ s$^{-1}$, or equivalently, a star formation surface density of $\sim 0.1 M_\odot$ yr$^{-1}$ kpc$^{-2}$, corresponds to superbubbles with a Mach number $\sim 5–10$. While the milder superbubbles can be produced by large OB associations, the latter kind requires super-starclusters. These derived conditions compare well with observations of disc galaxies with winds and the existence of multiphase halo gas. Furthermore, we find that contrary to the general belief that superbubbles fragment through Rayleigh–Taylor (RT) instability when they reach a vertical height of the order of the scaleheight, the superbubbles are first affected by thermal instability for typical disc parameters and that RT instability takes over when the shells reach a distance of approximately twice the scaleheight.

Key words: shock waves—ISM: bubbles—galaxies: ISM.

1 INTRODUCTION
Observations of nearby and high-redshift galaxies have shown that star formation in them often leads to galactic winds. Starburst galaxies, with star formation rate (SFR) in excess of a few tens of $M_\odot$ yr$^{-1}$ are known to excite such outflows. However, Heckman (2002) pointed out that it is not the average SFR, but the SFR surface density which is a deciding factor for the existence of outflows. He found a threshold SFR surface density of $\sim 0.1 M_\odot$ kpc$^{-2}$ yr$^{-1}$ as a pre-requisite for starbursts to be able to produce galactic winds.

The standard scenario of star formation leading to the wind phenomena posits that super-starclusters give rise to a large number of supernovae (SN) in a relative small region, which can produce a superbubble in the disc and can break out of the disc with enough momentum to produce a wind. Such super-star clusters, or young globular clusters, have been observed to have masses in the range of few $\times 10^5–6 \times 10^7 M_\odot$ within a typical radius of $\sim 3–10$ pc (Ho 1997; Martín-Hernández, Schaerer & Sauvage 2005; Walcher et al. 2005). The large amount of energy deposited into the interstellar medium (ISM) by these objects in the form of UV radiation and mechanical energy is believed to be an important feedback process. The mechanical energy from these super-starclusters has been shown to be important for the superbubble produced by the combined SNe to break out of the disc and produce a large-scale wind (e.g. Tenorio-Tagle, Silich & Muñoz-Tuñón 2003).

There have been a number of calculations, both analytical and numerical, dealing with the breakout of superbubbles from disc galaxies. The conditions for breakout depend strongly on the assumption of the stratification of gas in the disc. Consider an exponentially stratified disc with mid-plane ambient gas pressure $P_0$, gas density $\rho_0$, scaleheight $z_0$ and a bubble being blown by mechanical luminosity $L$. Mac Low & McCray (1988) defined a dimensionless parameter $D \equiv L \rho_0^{1/2} (P_0^{3/2} z_0)$, and noticed in their numerical simulations that break out of bubbles occurred when $D \geq 100$. The importance of this parameter can be understood by considering the self-similar evolution of a superbubble driven by an energy injection rate of $\mathcal{L}$, given by $r \sim (\mathcal{L} t^3 / \rho_0)^{1/5}$ and $\dot{r} \sim (3/5) (\mathcal{L} / \rho_0)^{1/5} t^{-3/5}$.

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This implies a speed of $\sim(3/5)(L/p_0 z_0^{3/2})^{1/3} \propto D^{1/3}$ when the superbubble reaches a distance of the scaleheight, for an ambient gas at a given temperature. According to this criterion, for a scaleheight $z_0 = 200$ pc, and mid-plane gas density $\rho_0 \sim 2.3 \times 10^{-24}$ g cm$^{-3}$, $P_0/\kappa_0 \sim n_0 10^4$ K cm$^{-3}$, a bubble with total mechanical luminosity of $\mathcal{L} \sim 3.8 \times 10^{37}$ erg s$^{-1}$ will be able to break out of the ISM.

Busa, Johnstone & Martin (1999) defined a dimensionless parameter $\beta = (27/154\pi)^{1/3} L_1^{1/3} p_0^{-1/3} \rho_0^{-1/3} z_0^{2/3}$ which is a ratio of the radius where the Mach number of the superbubble becomes unity, to the scaleheight. This is motivated by the self-similar solution of a stellar wind, $r \sim (125/154\pi)^{1/3} L_1^{1/3} p_0^{-1/3} \rho_0^{-1/3}$. They showed that this parameter is related to the above-mentioned $D$ parameter as $D = 17.9 b^2$. In other words, a superbubble with $b < 1$ is likely to be confined whereas blowout will occur for $b \geq 1$.

Koo & McKee (1992) analytically determined a condition for the breakout. Since the bubble accelerates after reaching a distance of the order of the scaleheight, owing to the rapidly decreasing density, it becomes liable to fragment due to Rayleigh–Taylor (RT) instability. If the Mach number of the bubble at scaleheight is $\geq 3$, then they argued that the bubble would be able to break out. They used radiative bubble model of Weaver et al. (1977) for a uniform density atmosphere in order to derive a critical mechanical luminosity for which the Mach number is unity, $L_{cr} \sim 17.9 p_0 z_0^2 c_s^3$, where $c_s$ is the isothermal sound speed of the ambient gas. The Mac Low & McCray condition of $D \geq 100$ translates to $L/L_{cr} \geq 5$. As we will find in our simulations, the Mach number of a bubble after breakout is of the order of $(1/5 c_s)(L/p_0 z_0^2)^{1/3}$. Therefore, the Mac Low–McCray condition of $D \geq 100$ translates to the condition that the Mach number at breakout is of the order of unity. We also note that they considered superbubbles that originated at a height from the mid-plane, which made it easier for bubbles to break out. Our simulations show that the critical luminosity for Mach number at a distance of the scaleheight to be unity is $L_{cr} \sim 125 p_0 z_0^2 c_s^3$, larger than the estimate of Koo & McKee (1992).

Koo & McKee (1992) then considered an additional strata of H II gas with a scaleheight of 1 kpc and mid-plane number density 0.025 cm$^{-3}$, and found the breakout condition to be of the order of $N_{OB} \sim 800$, or equivalently, $L \geq 4.1 \times 10^{38}$ erg s$^{-1}$. Sichl & Tenorio-Tagle (2001) considered the effect of halo gas pressure and determined a minimum energy for the superbubble to blow out of the galaxies (with both disc and spherical ISM distribution) with ISM gas mass in the range of $10^6$–$10^{10}$ M$_\odot$. For a disc galaxy with $M_{ISM} \sim 10^8$ M$_\odot$, they found a minimum energy of $\sim 10^{58}$ erg s$^{-1}$, corresponding to $N_{OB} \sim 100$.

As Heckman (2002) has emphasized, it is the surface density of SFR that determines the condition for the existence of galactic winds, and not the total luminosity. To translate the above energy conditions into a surface density, we need to estimate the surface area of such bubbles at the breakout epoch. In this paper, we re-visit this issue in order to understand the empirical threshold SFR surface density for galactic winds. Murray, Ménard & Thompson (2011) have recently argued that radiation pressure from UV radiation from a disc with an SFR surface density larger than 0.1 M$_\odot$ kpc$^{-2}$ yr$^{-1}$ can produce a large-scale wind. This estimate however crucially depends on the assumption of the grain opacity, and as Sharma & Nath (2012) have shown that the relevant opacity at UV may fall short of the requirements.

There have also been studies on the existence of multiphase gas in the haloes of spiral galaxies, and their connection to the star formation properties in the disc. Dahlem, Liskenfeld & Golla (1995) considered nine edge-on galaxies with extended synchrotron emitting halo gas, and derived a minimum value of surface density of energy injection for superbubble breakout, as $\sim 10^{-4}$ erg s$^{-1}$ cm$^{-2}$. Tüllmann et al. (2006) further considered X-ray, radio and far-infrared emission from the extended halo gas in a sample of 23 edge-on spiral galaxies, and found that the halo contained gas at low and high temperatures (multiphase) if the surface density of energy injection in the disc exceeds $\sim 10^{-3}$ erg s$^{-1}$ cm$^{-2}$. If the existence of multiphase halo gas depends on the process of superbubbles breaking out of the disc and depositing hot interior gas (as suggested by Tomisaka & Ikeuchi 1986; Tenorio-Tagle, Rozyczka & Bodenheimer 1990), as well as cold gas in the shell, then it would be interesting to compare the energetics of such superbubbles and the observed threshold energy injection rate.

In this paper, we study the standard scenario of thermal pressure of the gas interior to superbubbles being the driving mechanism for the wind, and derive a threshold condition for the superwind. We find that radiative loss of energy is important for the dynamics of shocks, and the inclusion of radiation loss increases the energy budget for the bubbles to break out of the disc and produce a wind. We also find that our estimate of the threshold energy requirement can explain the observed threshold SFR surface density for galactic outflows.

The paper is organized as follows. In Section 2, we derive an order-of-magnitude estimate of the threshold based on the key idea that the superbubble energetics needs to balance radiative cooling. Then we present the analytical formalism in Section 3 and discuss the results in Section 4. We then present the results from numerical simulations in Section 5, and discuss the effect of thermal and RT instability in Section 6.

### 2 Analytic Estimates

To begin with, we derive a threshold rate of SNe for a superbubble to continue to grow and ultimately break out of the disc from simple arguments. We can first consider the condition that the superbubble is able to drive a strong shock in the disc. This requires the volume energy injection time-scale to be shorter than the sound crossing time. In other words, if we consider a region of radius $R$ in the disc and an energy injection rate of $\mathcal{L}$, then one needs

$$\frac{1.5 n k T}{\mathcal{L} / (4\pi R^2/3)} \ll R/c_s, \quad (1)$$

where $c_s$ is the sound speed. This gives a lower limit of $\mathcal{L} / (\pi R^2) \gg 3 \times 10^{-6} n I_{4}\ E_{1/3}^1/2 \nu_{4}^{-1/3}$ erg s$^{-1}$ cm$^{-2}$, where $n$ is the ambient gas particle density in cm$^{-3}$ and $T = T_4 10^4$ K.

A second, and more stringent, constraint on SNe luminosity comes from accounting for radiative losses. Let us assume that when an SN remnant enters the radiative stage it quickly loses its energy and does not contribute to the energy input of the superbubble. Assume then that the radiative stage begins when the post-shock temperature is $T_s \simeq 2 \times 10^5$ K such that the radiation loss function is maximum and much larger than the minimum at $\sim 10^9$ K. We therefore define the time when an SN remnant loses its energy at time when the shock velocity is $v_s = 120$ km s$^{-1}$ (corresponding to the post-shock temperature of $2 \times 10^5$ K). It determines the corresponding time and radius as (see also Kahn 1998, who defined this as the beginning of phase III in the evolution of a bubble)

$$t_s = 1.4 \times 10^5 \frac{E_{1/3}}{n_{1/3}} \text{ yr}, \quad R_{\text{SN}} = 37 \frac{E_{1/3}}{n_{1/3}} \text{ pc}.$$  

(2)

One can therefore define the coherency condition as,

$$\frac{4\pi}{3} R_{\text{SN}}^3 t_s v_{\text{SN}} > 1,$$

(3)
which means that before an SN remnant stalls because of cooling losses, another SN explosion injects energy into the remnant and forms a single bubble. This condition determines the required SN rate

\[ v_{SN} > 30 \times 10^{-11} \left( \frac{n}{E_{51}} \right)^{4/3} \text{SN} \text{ yr}^{-1} \text{ pc}^{-3}. \]  

(4)

We can estimate the surface density of SNe, by multiplying this rate density by the scaleheight, which is the height of a bubble at the epoch of breakout. For a scaleheight of 500 \( z_0 = 0.5 \) pc, this corresponds to \( 1.5 \times 10^{-2} (n/E_{51})^{4/3} z_0 \text{SN} \text{ yr}^{-1} \text{ kpc}^{-2} \). (The scaleheight is relevant here because, as we will see later, the maximum radius of bubbles in the plane parallel to the disc is of the order of \( \pi z_0 \).) Finally, we recall that for a Salpeter IMF, one SN corresponds to 150 M_\odot of stellar mass, considering stars in the range of 1–100 M_\odot. Therefore, the threshold condition for SFR surface density becomes \( \sim 2.5(n/E_{51})^{4/3} z_0 \text{ SN} \text{ yr}^{-1} \text{ kpc}^{-2} \). The corresponding surface density of energy injection is \( \sim 0.05 n^{1/3} E_{51}^{-1/3} z_0 \text{ erg s}^{-1} \text{ cm}^{-2} \). It is interesting to find that these above estimates of the threshold energy injection or SFR surface density are comparable to the observed threshold for the existence of multiphase halo gas (Tüllmann et al. 2006) and superwinds (Heckman 2002).

3 KOMPANEETS APPROXIMATION

We first discuss the expansion of blastwaves in a stratified atmosphere, in the adiabatic case and then for radiative shocks. Kompansets (1960) had first analytically worked out the case of adiabatic shocks in this case (see e.g. Bisnovatyi-Kogan & Silich 1995). Consider an exponentially stratified medium described by \( \rho(z) = \rho_0 \exp(-z/z_0) \), where \( \rho_0 \) is the mid-plane density and \( z_0 \) is the scaleheight and \( E_0 \) is the explosion energy. It is assumed that the shock pressure is uniform, and is given by

\[ P_{sh} = \frac{(y - 1)\lambda E_0}{V}, \]  

(5)

where \( \lambda \sim 1 \) (Kompaneets 1960) is a constant that differentiates the shock pressure from the average pressure inside the bubble; Bisnovatyi-Kogan & Silich (1995) evaluated \( \lambda = 1.33 \). We use \( \lambda = 1 \) for simplicity. We define a dimensionless time-like parameter as

\[ y = \int_0^t \sqrt{\frac{(y^2 - 1)E_0}{2\rho_0 V}} \, dt. \]  

(6)

Where \( E_{sh} \) is the thermal energy of the interior gas, \( V \) is the volume of the bubble and \( t \) is the time. The shape of the shock front is derived as

\[ r = 2z_0 \arccos \left\{ \frac{1}{2} \exp(z/2z_0) \left[ 1 - \frac{y^2}{4z_0^2} + \exp(-z/z_0) \right] \right\}. \]  

(7)

The location of the top and bottom of the bubble then follows by setting \( r = 0 \) with \( y = y(z_0) \),

\[ z_\pm(y) = -2z_0 \ln(1 \mp y/2), \]  

(8)

which shows that the top of the bubble reaches infinity when \( y \to 2z_0 \) while \( t \) remains finite. This implies that the bubble accelerates in the \( z \)-direction due to stratification, after an initial deceleration phase when the bubble is small and spherical, as in the usual Sedov–Taylor solution. The maximum cylindrical radius of the bubble is also obtained from the above solution by putting \( \partial r/\partial z = 0 \),

\[ r_{max}(\hat{y}) = 2z_0 \arcsin(\hat{y}/2). \]  

(9)

The \( z \)-component of the velocity of the topmost point of the bubble is given by

\[ v_z(\hat{y}) = \frac{1}{1 - \hat{y}^2/2} \sqrt{\frac{(y^2 - 1)E_0}{2\rho_0 V(t)}}. \]  

(10)

3.1 Continuous energy injection

We can extend Kompaneets approximation and radiative blastwave calculation to the case of continuous energy injection. Schiano (1985) performed a similar calculation in the case of an active galactic nucleus. Consider an association with \( N_{OB} \) stars with masses above 8 M_\odot, which ultimately produce supernovae. If we consider the main-sequence lifetime as \( \tau_{SN} \sim 5 \times 10^7 \) yr for these stars, then the total mechanical luminosity of the SN in the association can be written as,

\[ \mathcal{L} = 6.3 \times 10^{33} N_{OB} E_{51} (\tau_{SN}/5 \times 10^7 \text{ yr})^{-1} \text{ erg s}^{-1}, \]  

(11)

where supernova energy is \( 10^{51} E_{51} \) erg. As McCray & Kafatos (1987) have argued, since the main-sequence lifetime scales with \( M \) as \( \propto M^{-1.6} \), and since the initial mass function (IMF) is given by \( dN/M \propto M^{-1.35} \), for a Salpeter IMF, the rate of SN will scale with time as \( \propto t^{1.35/1.6 - 1} \propto t^{1.35/1.6 - 1} \), which is roughly constant in time. Here we have used \( \frac{dN}{dV} \propto t^{1.35/1.6 - 1} \), given the above-mentioned dependence of stellar main-sequence lifetime. Therefore, we can write, for the adiabatic case, the total energy in the superbubble as \( E_{sh} = \mathcal{L} t \).

Instead of equation (10), the \( z \) velocity of the top of the bubble is then given by

\[ v_z(\hat{y}) = \frac{1}{1 - \hat{y}^2/2} \sqrt{\frac{(y^2 - 1)E_0}{2\rho_0 V(t)}}. \]  

(12)

and the corresponding \( y \) parameter is also written in terms of \( t \), as

\[ y = \int_0^t \sqrt{\frac{(y^2 - 1)\mathcal{L}}{2\rho_0 V(t')}} \, dt'. \]  

(13)

These equations can determine the dynamics of the superbubble in the case of continuous energy injection.

3.2 Radiative loss with continuous injection

Radiative losses can be important for the dynamics of both the blastwave and a superbubble with continuous energy injection. Shocks become radiative when the cooling time \( t_{cool} \ll t \). The cooling time behind the shell can be estimated as \( t_{cool} = 1.5T/(4n(T)\lambda) \), for a strong shock with \( n = n_0 \exp(-z/z_0) \) and the shock temperature being estimated from the shock speed (in the \( z \)-direction, say). We assume a cooling function, as given by equation 12 in Sharma, Quataert & Parrish 2010, appropriate for gas with \( n(T) = 3.57 \times 10^{-6} T_{keV}^{1.7} + 5.8 \times 10^{-2} T_{keV}^{0.5} + 6.3 \times 10^{-2} \) erg s \(-1\) cm \(^3\).\( T > 0.02 \) keV,

\[ = 6.72 \times 10^{-2}(T_{keV}/0.02)^{0.6} \text{ erg s}^{-1} \text{ cm}^{-3}, \]  

\[ T \leq 0.02 \text{ keV}, \quad T \geq 0.0017 \text{ keV} \]

\[ = 1.544 \times 10^{-2}(T_{keV}/0.0017)^{0.6} \text{ erg s}^{-1} \text{ cm}^{-3}, \]  

\[ T < 0.0017 \text{ keV} , \]  

(14)
where $T_{\text{keV}}$ is the temperature in keV. Fig. 1 shows the ratio $t_{\text{cool}}/t$ as a function of the bubble height $z_+$ for bubbles with continuous energy injection for a few cases. The curves show that the shock enters the radiative phase much before reaching the scale-height unless the ambient density and scaleheight are very small and $N_{\text{OB}}$ is very large (e.g. the case with $n_0 = 0.1 \text{ cm}^{-3}$, $z_0 = 200 \text{ pc}$, $N_{\text{OB}} = 5000$).

Radiation loss from the shocked medium can therefore be important (see also Maciejewski & Cox 1999). Kovalenko & Shchekinov (1985) had calculated the dynamics of a blastwave with radiative loss, assuming that the shock kinetic energy is converted into thermal energy of gas in a thin shell behind it, and that radiative loss from this shell keeps the shock isothermal. It can then be shown that for a strong shock the energy lost per unit mass is $\sim (1/2)n_0 c_s^2$, where $c_s$ is the shock speed. From the Hugoniot condition for a strong shock

$$u_s^2 = (y^2 + 1)P_s^{\gamma} = (y^2 - 1)E_{\text{th}}^{\gamma}$$

where $E_{\text{th}}$ is the thermal energy of the shocked gas. The structure of the shock in this case is such that the interior gas remains hot and adiabatic, whereas the shocked ambient gas that is swept into a shell loses its energy radiatively and is kept at a constant temperature (at $\sim 10^7 \text{ K}$). We note that Mac Low & McCray (1988) showed that the radiative loss from the interior hot gas of the bubble does not change the dynamics of the bubble.

Following the calculation of Kovalenko & Shchekinov (1985) for a radiative blastwave, we assume that bubbles with continuous energy injection also form an isothermal thin shell, after a certain time $t_1$ when it enters the radiative phase. For simplicity, we also assume a self-similar solution for a spherical shock, of the type given by Weaver et al. (1977), $r_s = A t^{1/2} r^{3/2}$, where $A$ is a constant depending on the ambient density. Furthermore, Weaver et al. (1977) have pointed out that a fraction 6/11 of the total energy is stored in the shell and the rest in the rarefied gas inside. In the spirit of Kovalenko & Shchekinov (1985), we assume the total shell energy to be thermal in nature. In other words, initially $E_{\text{th}} = (6/11)Ct$.

We can determine the time evolution of $E_{\text{th}}$ as follows.

Using the result derived in equation (15) that the amount of energy lost per unit volume is $(1/2)\rho_0 u_s^2 = (y^2 - 1)E_{\text{th}}/(4V(r))$, we can write for the evolution of thermal energy in this case,

$$E_{\text{th}}(r) = \frac{6}{11}Ct - \pi (y^2 - 1) \int_{r_1}^r \frac{E_{\text{th}}(r)}{V} r^2 \, dr$$

$$= \frac{6}{11} C \left( \frac{r}{A} \right)^{5/3} - \pi (y^2 - 1) \int_{r_1}^r \frac{E_{\text{th}}(r)}{V} r^2 \, dr$$

(16)

Here $r_1$ is the radius at time $t_1$. We can explicitly solve this equation for a spherical shock, and then use the results to estimate the $z$-velocity of an oval shaped bubble. For a spherical shock (with volume $V = \frac{4}{3}\pi r^3$), the energy equation (16) can be shown to yield a solution of the type $E_{\text{th}}(r) = br^\alpha$, where

$$b \left[ \alpha + \frac{3}{5} (y^2 - 1) \right] r^{\alpha - 1} = \frac{6}{11} \frac{5}{3} \frac{C}{A^{5/3}} r^{2/3}$$

(17)

Comparing the powers of $r$ from both the sides we get, $\alpha = \frac{5}{3}$. Putting this value of $\alpha$ in equation (17) and comparing the coefficients of time on both sides we get,

$$b = \frac{30}{99} \frac{C}{A^{5/3}}$$

(18)

Therefore $E_{\text{th}}(r)$ becomes,

$$E_{\text{th}}(r) = 0.3 Ct$$

(19)

showing that roughly 70 per cent of the total energy is radiated away.

Note that this is an asymptotic value of the loss in the limit $r \gg r_1$, in the regime where the approximation $E \propto r^5$ is valid. We can therefore use equations (12) and (13), with the above value of $E_{\text{th}}$, and determine the dynamics of a radiative superbubble with continuous energy injection.

### 4 Analytic Results

Fig. 2 shows the evolution of the Mach number for a $10^4 \text{ K}$ gas as a function of time, for an adiabatic blastwave, a superbubble with continuous energy injection with and without radiative loss. It is convenient to define a dynamical time-scale for this problem (Mac Low & McCray 1988), $t_L \sim z_0^{4/3} (\rho_0/L)^{1/3}$, which is the

![Figure 2](https://example.com/figure2.png)

**Figure 2.** The evolution of the ratio of $v_1$ to $c_s$ (the sound speed for an ambient gas at $10^4 \text{ K}$) is plotted against time, for an adiabatic blastwave (thick solid line), adiabatic superbubble with continuous energy injection (dashed) and with radiative loss (dotted line).
expected time to reach the scaleheight for a self-similar evolution of superbubbles. For $z_0 = 200 \text{ pc}$, $L \sim 1.3 \times 10^{37} \text{ erg} \text{ s}^{-1}$ and $\rho_o \sim 10^{-25} \text{ g} \text{ cm}^{-3}$ (for $\mu \sim 0.6$), $t_0 \sim 2.8 \text{ Myr}$. We find that the $z$-velocity shows a minimum at $\sim 1.5t_0$, when it reaches a distance of the scaleheight. We denote this minimum value of $z$-velocity as $v_{z, \text{min}}$, and refer to this epoch as the ‘stalling epoch’ in our discussion below.

Fig. 3 shows the Mach number at stalling height, as a function of $L$, the mechanical luminosity (which scales as $N_{\text{OB}}$). Interestingly, superbubbles with Mach number (at stalling height) of order less than unity can be triggered by even a single SN. These, in principle, can accelerate later and therefore break out of the disc. However, as we shall see later with our simulations, there is a minimum number of SNe needed for superbubbles to break out of the disc, particularly for high-density discs. We also find from Fig. 3 that in order to achieve a Mach number at stalling height of the order of $\sim 5$, one needs $L \geq 7 \times 10^{38} \text{ erg} \text{ s}^{-1}$, for $n_0 = 1 \text{ cm}^{-3}$ and $z_0 = 500 \text{ pc}$. This is larger than the estimate of Koo & McKee (1992) and Mac Low & McCray (1988), because of the inclusion of radiative loss from the shell. If we consider $v_{z, \text{min}}/c_s \geq 5$ as the breakout condition, then we find that larger densities and scaleheights put more stringent condition on the bubble to break out.

Next, we plot in Fig. 4 the minimum Mach number as a function of the surface density of $N_{\text{OB}}$, considering the surface area of the bubble at the stalling height. Note that we are not concerned with the mean surface density of SFR in the disc galaxy here. The energy injection considered here is localized, but the relevant surface area as far as an emerging superbubble is concerned, is the area of the bubble in the plane of the disc at the point of breaking out. We find that for the surface density of energy deposition the analytic curves become independent of the scaleheight and depend only on the gas density and number of SNe. This is because the area of a superbubble in the plane parallel to the disc, scales with $z_0^2$, and is a constant for a given scaleheight. We find that for a scaleheight of $500 \text{ pc}$, the threshold surface density of SNe is $N_{\text{OB}} \sim 1000 \text{ kpc}^{-2}$.

5 NUMERICAL SIMULATIONS

In addition to analytic estimates and approximate calculations, we have performed 2-D axisymmetric hydrodynamic simulations of breakout using the ZEUS-MP code (Hayes et al. 2006). ZEUS-MP is a publicly available, second-order accurate Eulerian hydrodynamics code. We have carried out two sets of simulations: the first set compares numerical simulations with the analytic Kompaneets calculation of strong shocks in stratified atmospheres (hereafter these runs are referred to as ‘Kompaneets runs’), and the second set of calculations use a more realistic setup, such as disc gravity, mass loading of the ejecta, for shock (superbubble) breakout in star-forming galaxies (hereafter these runs will be called ‘realistic runs’).

In this section, we introduce the equations that we solve numerically, the initial and boundary conditions, and the choice of setup parameters. The simulations are run using the 2-D axisymmetric, spherical polar ($r, \theta, \phi$) coordinates.

5.1 Governing equations

We solve the following standard Euler’s hydrodynamic equations including cooling, external gravity, and mass and energy loading at inner radii:

$$\frac{d\rho}{dr} = -\rho \nabla \cdot \mathbf{v} + S_p(r),$$

$$\frac{d\mathbf{v}}{dr} = -\nabla p + \rho \mathbf{g},$$

$$\frac{de}{dr} = -q^-(n, T) + S_e(r),$$

where $d/dr \equiv \partial/\partial r + \mathbf{v} \cdot \nabla$ is the Lagrangian derivative, $\rho$ is the mass density, $\mathbf{v}$ is the fluid velocity, $p$ is the thermal pressure, $e = p/(\gamma - 1)$ is the internal energy density (we use $\gamma = 5/3$ valid for an ideal non-relativistic gas), $\mathbf{g} = -\text{sgn}(z)g \hat{z}$ (sgn[z] = ±1 for $z \geq 0$) is the constant external gravity pointing towards the $z = 0$ plane, $q^- = n_e n_i \Lambda(T)$ is the cooling term due to radiation where $n_e$ and $n_i$ are the electron and ion number densities, $\Lambda(T)$ is the cooling function (as given in equation 14). There are source terms in the mass and internal energy equations ($S_m, S_e$). These terms are non-zero and constant only within $r_{\text{in}}$, a small
We first describe the results of our Kompaneets runs, of superbubbles in a stratified atmosphere without external gravity or mass loading. Fig. 4 shows the variation of the minimum Mach number of the top of the superbubble as a function of the surface density of energy injection in the disc, for a scaleheight of 200 pc and two values of ambient density, $n_0 = 0.1$ and $1 \text{ cm}^{-3}$. We find that the analytical results overestimate the Mach number of the superbubbles compared to the simulations by a factor of the order of $10^4$ for the case of large ambient density ($1 \text{ cm}^{-3}$), because the analytical estimate of energy loss described in the previous section is based on simplified assumptions. Note that since we determine the value

5.3 Kompaneets runs

Table 1. Parameters for Kompaneets runs ($\mathcal{L} = 6.3 \times 10^{35} \text{erg s}^{-1} N_{\text{OB}}$).

| $n_0$ (cm$^{-3}$) | $N_{\text{OB}}$ | $r_{\text{in}}$ (pc) | $r_{\text{min}}$ (pc) | $r_{\text{max}}$ (pc) |
|------------------|---------------|---------------------|---------------------|---------------------|
| 0.1              | 1             | 10                  | 5                   | 2500                |
| 0.1              | 10            | 21                  | 10                  | 2500                |
| 0.1              | 100           | 44                  | 30                  | 3000                |
| 0.1              | 300           | 63                  | 40                  | 3000                |
| 0.1              | 1000          | 94                  | 70                  | 3500                |
| 1.0              | 1             | 5                   | 3                   | 2500                |
| 1.0              | 10            | 10                  | 5                   | 2500                |
| 1.0              | 100           | 20                  | 10                  | 2500                |
| 1.0              | 300           | 29                  | 15                  | 3000                |
| 1.0              | 1000          | 44                  | 30                  | 3500                |

Table 2. Parameters for realistic runs.

| $z_0$ (pc) | $n_0$ (cm$^{-3}$) | $N_{\text{OB}}$ | $r_{\text{in}}$ (pc) | $r_{\text{min}}$ (pc) | $r_{\text{max}}$ (pc) |
|-----------|------------------|---------------|---------------------|---------------------|---------------------|
| 100       | 0.1              | 1             | 10                  | 5                   | 1000                |
| 100       | 0.1              | 10            | 21                  | 10                  | 2500                |
| 100       | 0.1              | 100           | 44                  | 10                  | 2500                |
| 100       | 0.1              | 300           | 63                  | 10                  | 2500                |
| 100       | 0.1              | 1000          | 94                  | 50                  | 2500                |
| 100       | 1               | 100           | 20                  | 10                  | 2500                |
| 100       | 1               | 300           | 29                  | 15                  | 2500                |
| 100       | 1               | 1000          | 44                  | 30                  | 2500                |
| 100       | 1               | 2000          | 55                  | 40                  | 2500                |
| 100       | 1               | 3000          | 63                  | 40                  | 2500                |
| 500       | 0.1              | 10             | 21                  | 10                  | 2500                |
| 500       | 0.1              | 100            | 44                  | 10                  | 2500                |
| 500       | 0.1              | 300            | 63                  | 30                  | 3500                |
| 500       | 0.1              | 1000           | 94                  | 50                  | 3500                |
| 500       | 0.1              | 3000           | 135                 | 110                 | 3500                |
| 500       | 0.1              | 10 000         | 201                 | 160                 | 3500                |
| 500       | 0.1              | 50 000         | 344                 | 300                 | 12 000              |
| 500       | 0.1              | 100 000        | 433                 | 400                 | 12 000              |
| 500       | 1               | 1000           | 44                  | 30                  | 2500                |
| 500       | 1               | 2000           | 55                  | 40                  | 2500                |
| 500       | 1               | 3000           | 63                  | 40                  | 2500                |
| 500       | 1               | 5000           | 75                  | 50                  | 3500                |
| 500       | 1               | 10 000         | 94                  | 70                  | 3500                |
| 500       | 1               | 100 000        | 201                 | 150                 | 5500                |
of \( z_0 \) by the position of the maximum density, clumps in the shell formed due to thermal instability (see below for details) introduce some uncertainty. This manifests in the kinks seen in the simulation results in Fig. 4 and also later in Fig. 6.

5.4 Realistic runs

Next, we describe simulations that includes vertical disc gravity and mass loading. We study the case of ambient gas at \( T = 10^4 \) K, with mid-plane densities \( n_0 = 0.1 \) and \( 1 \) cm\(^{-3}\), and scaleheights \( z_0 = 100 \) and 500 pc.

Our choice of parameters essentially brackets the possible range of gas density and scaleheight in disc galaxies. For example, the distribution of the extraplanar gas in Milky Way has two components, that of warm ionized gas and cold H\( \alpha \). The warm ionized gas has been observed to have an exponential profile with \( n_0 \sim 0.01\sim 0.03 \) cm\(^{-3}\) and \( z_0 \sim 400\sim 1000 \) pc (Reynolds 1991; Nordgren, Cordes & Terzian 1992; Gaensler et al. 2008). For H \( \alpha \) distribution, Dickey & Lockman (1990) found that the vertical distribution is best described by a Gaussian with full width at half-maximum of 230 pc and a central density of 0.57 cm\(^{-3}\). The combined distribution of these two components are bracketed by exponentials with the scaleheights and mid-plane densities assumed here.

We also use smaller scaleheights in our simulations. The scaleheight near the centres of galaxies is smaller than that in the outer regions, because of deeper gravitational potentials in the central regions. Also Dalcanton, Youchim & Bernstein (2004) found that the H\( \alpha \) scaleheight of disc galaxies varies with the rotation speed (or, equivalently, the galactic mass). Dwarf spirals with rotation speed \( \sim 50 \) km s\(^{-1}\) have \( z_0 \sim 200 \) pc, whereas larger galaxies (with rotation speed in excess of 120 km s\(^{-1}\)) have \( z_0 = 500\sim 1000 \) pc. Also, as Basu et al. (1999) have found, the scaleheight encountered by Milky Way superbubbles such as W4 is rather small (\( \sim 100 \) pc).

We first find that unlike in the analytical case, where superbubbles ultimately break out of the disc sooner or later, irrespective of the energetics, the realistic simulation runs show that for high-density disc material \( (n_0 \geq 1 \) cm\(^{-3}\)), superbubbles keep decelerating for ever for a surface density of OB stars \( \sim 100(z_0/100 \) pc) kpc\(^{-2}\). In other words, superbubbles never break out of the disc in these cases. The corresponding energy injection surface density is \( \sim 2\sim 5 \times 10^{-2} \) erg cm\(^{-2}\) s\(^{-1}\). For lower density ambient gas, \( n_0 \sim 0.1 \) cm\(^{-3}\), however, even a single SN event can drive a bubble through the disc. We note that this limit is consistent with that found by Silich & Tenorio-Tagle (2001) for a Milky Way type disc.

In the case of a superbubble breaking out of the disc, there are differences in the way they evolve depending on the energy injection rate. We show the evolution of the speed of the topmost point of the bubble as a function of time for four cases in Fig. 5, for two mid-plane densities \( n_0 = 0.1 \) and \( 1 \) cm\(^{-3}\) and two scaleheights \( z_0 = 100 \) and 500 pc, all for a surface density of OB stars of 1000 kpc\(^{-2}\). The curves show that the bubbles show acceleration after breaking out of the disc only for the case of low density and small scaleheight (see the curve at the top-left corner, for \( n_0 = 0.1 \) cm\(^{-3}\), \( z_0 = 100 \) pc). In other cases, for disc column density \( \geq 3 \times 10^{19} \) cm\(^{-2}\), the bubbles either coast along with the speed that they reach at the breakout, or decelerate to some extent, for a considerable period of time before they start accelerating after reaching a distance of several scaleheights. The curves show that the speed at the stalling height, or the minimum speed of the bubbles, is an important characteristics of the bubble dynamics. It is important because this is the characteristic speed with which the bubble sweeps most of the extra-planar region of the halo. Also, since the bubble begins to accelerate only after reaching a distance of a few times the scaleheight, the corresponding RT instability should not set in at the scaleheight, but at a much larger distance. We shall re-visit this point in the next section on instabilities. In some cases, the curves show a deceleration at late times. This is due to the formation of clumps in the shell from radiative cooling, which often sink through the hot gas owing to gravity.

We have found that typically the minimum speed \( v_{\text{min}} \sim (1/5)\left(L/\rho_{\text{0}}z_0^2\right)^{1/3} \sim z_0/(5t_0) \), where \( t_0 = \text{the dynamical time defined earlier. These values are shown as horizontal lines in Fig. 5 for respective cases. It is easy to see that in the case of little radiation loss, the speed of the bubble at the time of reaching the scaleheight is \( \sim (3/5)\left(L/\rho_{\text{0}}z_0^2\right)^{1/3} \), as expected from the self-similar evolution of a bubble \( (r \sim \left(Lt^3/\rho_{\text{0}}\right)^{1/3}) \). Our simulations show that the actual speed is roughly a third of this value, and therefore shows the importance of radiative loss in the dynamics of superbubbles. As analytically derived earlier, radiation losses remove as much as 70 per cent of the total energy of the superbubbles. We recall that for an ambient medium with a given temperature, the dimensionless quantity defined by Mac Low & McCray (1988) is \( D = (5v_{\text{min}}/c_{\text{s}})^3 \), so that their condition of \( D \geq 100 \) for break out corresponds to a minimum Mach number of the order of unity.

We show the resulting value of minimum Mach number of superbubbles for different \( n_0 \) and \( z_0 \) in Fig. 6, as a function of surface density of energy injection. The curves show that in terms of energy injection or SNe surface density, the crucial parameter is the mid-plane gas density, which separates the curves, as was also indicated by our analytical results. Superbubbles with a given surface density of energy injection find it easier to break out of discs with lower mid-plane density. However, scaleheight also makes a small difference unlike in the analytical calculations; a higher energy density is required to clear a thicker disc.

The important features of our results as shown in Fig. 6 are as follows.

(i) As mentioned above, the condition for a breakout from a dense ambient medium with gas density of \( n_0 = 1 \) cm\(^{-3}\) is an energy injection rate surface density of \( 2\sim 5 \times 10^{-3} \) erg cm\(^{-2}\) s\(^{-1}\). For lower gas densities, the required rate density is \( \sim 10^{-6} \) erg cm\(^{-2}\) s\(^{-1}\). The
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and the corresponding requirement on SFR surface density increasing to \( \sim 0.1 \, \text{M}_\odot \, \text{yr}^{-1} \, \text{kpc}^{-2} \), the observed threshold. Therefore, the Heckman (2002) threshold (\( \sim 0.1 \, \text{M}_\odot \, \text{yr}^{-1} \, \text{kpc}^{-2} \)) for superwinds corresponds to a larger requirement on the part of superbubbles, of not only breaking out of discs but doing so with a large Mach number.

6 THERMAL AND RAYLEIGH–TAYLOR INSTABILITY

The focus till now was on the important \( v_\parallel/c_s \) parameter (the minimum Mach number of the shell) which determines the fate of the superbubble after it crosses the scaleheight. In this section, we discuss the role of different instabilities, in particular RT and thermal instabilities, in our 2-D breakout simulations. When the superbubble reaches a scaleheight, the shock is generally believed to accelerate owing to the decrease in pressure. This should lead to the onset of the RT instability, as has been invoked in previous analytical works (e.g. Koo & McKee 1992) and seen in numerical simulations (e.g. Mac Low, McCray & Norman 1989). However, as mentioned earlier, our simulations show that superbubbles do not accelerate until after they reach a distance of several scaleheights (as was also suggested by Ferrara & Tolstoy 2000 who assumed spherical bubbles). Therefore, RT instability occurs at a distance much larger than the scaleheight. Also we find that before the onset of RT instability, the superbubble expanding in the disc suffers from thermal instability in the early stages of its evolution. This instability leads to clumping and fragmentation of the shell of the superbubble well in advance of the RT instability, and can therefore affect the outcome of the RT instability.

Fig. 7 shows the 2-D snapshots of temperature at two different times for our fiducial high-resolution run (\( N_{\text{OB}} = 5000, n_0 = 1 \, \text{cm}^{-3}, z_0 = 500 \, \text{pc} \)). Fig. 6 indicates that the minimum Mach number for this case is \( \approx 2 \) and the bubble is just about able to break out within the starburst time-scale. The temperature snapshot at early time (9 Myr), when the bubble has just reached the scaleheight, shows that the bubble is roughly spherical. The radiative shell seems to develop corrugations where the hot bubble gas and the radiatively cooled shocked gas interpenetrate. The shell is at \( \approx 10^7 \, \text{K} \) (the same as the ambient ISM temperature), the temperature below which the cooling function drops suddenly and the gas becomes thermally stable. The dense shell is more clearly seen in the density snapshots of Fig. 8. The corrugations are definitely driven by radiative cooling.

Figure 6. The minimum Mach number of the top of the bubble in our realistic runs are shown as a function of \( N_{\text{OB}} \) per kpc\(^{-2} \), and \( L/\pi r^2 \) (erg cm\(^{-2} \) s\(^{-1} \)) for \( n_0 = 0.1, 1 \, \text{cm}^{-3} \) and \( z_0 = 100, 500 \, \text{pc} \). Note that, for \( n_0 = 1 \, \text{cm}^{-3} \), the shocks stall for a surface density of OB stars \( \leq 500 \, \text{kpc}^{-2} \). The cases for which \( t_{\text{cool}} < t_f \), are shown by darkened points, these cases are marked by thermal instability.

Figure 7. Temperature contours (colour coded) for a superbubble with \( N_{\text{OB}} = 5000, n_0 = 1 \, \text{cm}^{-3}, z_0 = 500 \, \text{pc} \), at \( t = 9 \, \text{Myr} \), when the top of the bubble has reached a distance of the scaleheight (left-hand panel), at 39.3 Myr, when it has reached a distance \( \sim 3z_0 \) (middle panel). The rightmost panel shows the case of the same superbubble without radiative cooling at \( t = 39.3 \, \text{Myr} \), the same evolutionary epoch as the middle panel.
because the run without radiative cooling shows a smooth shell (the third panel in Figs 7 and 8).

While the fragments of cold shell are confined to the bubble boundary at early times, the cold gas lags behind the hot gas at later times because the hot gas is pushed out by supernova heating. The cold blobs are only pushed out because of the drag force due to the hot gas but eventually trail behind. The cold blobs embedded in the hot gas are reminiscent of the cold multiphase filaments observed in galactic outflows, such as M82. Since in our simulations cold gas leaves the simulation box from the inner boundary, all the cold blobs embedded in the hot bubble come from the fragmenting cold shell. In reality, some cold gas from the cold star-forming regions can also be uplifted by the hot gas. At late times, in the runs with cooling, there are some signs of bubble breaking out because of RT instability close to the polar regions. All such signatures of RT instability are missing in the run without cooling (panel 3). This is mainly because RT instability in the run with cooling is seeded with large amplitude perturbations by corrugations caused by shell cooling.

In order to assess the relative importance of thermal and RT instabilities, we compare the two time scales in Fig. 9. We note that the time-scale for RT instability \( t_{\text{RT}} = \sqrt{\frac{T_0}{v_z + g}} \) is comparable to the free-fall time \( t_{\text{ff}} = \sqrt{\frac{2z}{v_z + g}} \), for the largest mode with \( k \sim 2\pi/\lambda \), where \( z \) and \( v_z \) are the height and acceleration of the shell, and \( g \) is the acceleration due to gravity. We plot this time-scale with a solid line in Fig. 9, along with the cooling time \( t_{\text{cool}} = 1.5kT_0/\Lambda \) of the shell as a function of time for runs corresponding to Figs 7 and 8. We use the position of the outermost densest part to identify the shell position. In the left-hand panel of Fig. 9, we show the case of \( N_{\text{OB}} = 5000, n_0 = 1 \text{ cm}^{-3}, z_0 = 500 \text{ pc} \). We expect the shell to cool radiatively if \( t_{\text{cool}} \) is shorter than time. And indeed, the radiative cooling time is shorter than time at early times. This is consistent with the cooling and fragmentation of the dense shell seen in Figs 7 and 8. One point of caution: we should ideally plot the cooling time of the shell assuming the shell temperature and density corresponding to an adiabatic shock because cooling will happen if this time-scale is short. Here we are plotting the cooling time of the shell, which for the left-hand panel case, has already cooled to low temperatures. Since cooling time increases sharply below \( 10^5 \text{ K} \), \( t_{\text{cool}} \) is barely smaller than time in the left-hand panel of Fig. 9. At later times \( t_{\text{cool}} \) becomes longer than time and we do not expect the newly accumulated shell material to cool. The RT time-scale \( (=t_{\text{ff}}) \) is always longer than time for the fiducial run. The free-fall time increases initially as the shock slows down until a scale-height. After that the shock moves at a small Mach number \( \sim 2 \). This is consistent with the fact that we do not see vigorous RT instability in Figs 7 and 8.

The middle panel of the Fig. 9 shows various time-scales for a mid-plane density of \( n_0 = 0.1 \text{ cm}^{-3} \). The cooling time for this case is shorter than the cooling time for the higher density case. This seems inconceivable given the higher density and efficient cooling for the run in the left-hand panel. This discrepancy arises because although the density for the \( n_0 = 0.1 \text{ cm}^{-3} \) is smaller, the temperature of the post-shock gas is \( 10^3 \text{ K} \), where the cooling function peaks. Consequently the cooling time is shorter than the higher density run. For comparison, we have also plotted the cooling and free-fall time-scales for the runs without cooling in the right-hand panel. The density and temperature snapshots for this run do not show cooling-induced fragmentation.

We have shown in Fig. 6 by darkened points the cases in which \( t_{\text{cool}} \) is always less than \( t_{\text{ff}} \) for different values of \( \frac{L}{\pi \rho^2 z_0} \), and \( n_0 \). We find that these cases mostly appear for which, roughly, \( 10 \geq v_z, \text{min} \geq 2 \), except for the case of \( z_0 = 500 \text{ pc} \) and \( n_0 = 1 \text{ cm}^{-3} \), for which there is a cross-over point in time after \( t_{\text{cool}} \geq t_{\text{ff}} \). We note that this range of \( v_z, \text{min} \) corresponds to a case in which the shell temperature \( T_s \) remains in the range of \( 2 \times 10^4 \leq T_s \leq 10^5 \), where the cooling function peaks. This implies a range in \( N_{\text{OB}} \) for which thermal instability is important. In the low \( N_{\text{OB}} \) limit, the shock is not strong enough and \( T_s \leq 10^4 \), and in the high \( N_{\text{OB}} \) case, the shock is very strong \( (T_s > 10^5 \text{ K}) \) and \( t_{\text{ff}} \) (RT time-scale) is shorter than \( t_{\text{cool}} \) at late times.

We are therefore led to conclude that superbubbles are affected not only by RT instability but also by thermal instability, depending on the density and energy injection. This implies that the fragmentation of the bubble shell that releases the hot interior gas into the halo occurs under the combined effects of thermal instability at early times and RT instability at late times if the Mach number at stalling epoch is large enough.

7 DISCUSSION AND SUMMARY

Superbubbles with fragmented shells are believed to ultimately form ‘chimneys’ (Norman & Ikeuchi (1989), which connect the halo gas to the processes in the disc in different ways. Apart from transporting hot gas to the halo, chimneys provide a natural channel...
for Lyman continuum photons from hot stars in the disc to reach
the diffuse ionized medium of the Reynolds layer (Reynolds 1991; Døve & Shull 1994). It is however important for the superbubble
shells to fragment before the main-sequence lifetimes of O stars for
a substantial fraction of ionizing radiation to escape the disc (Døve,
Shull & Ferrara 2000). This implies a fragmentation time-scale of
∼3–5 Myr, which is comparable to the dynamical time-scale
(τa ∼ 5/3 ⋅ ρ0 ⋅ L1/3), for superbubbles with L ∼ 1038 erg (corre-
ponding to NOB ∼ 200), typical disc parameters. This is the energy
scale for the largest of the OB associations, and as our results show
superbubbles with smaller energetics find it hard to pierce through
the disc, unless the OB association is located much above the mid-
plane level.

In other words, for superbubbles to act as effective conduits of
ionizing radiation for the halo, or for the intergalactic medium (at
high redshift, in the context of the epoch of reionization), the su-
perbubbles need to fragment roughly around the time when they
reach a scaleheight. This is unlikely to happen only through RT in-
stability as superbubbles do not accelerate until reaching a distance
of several scaleheights. Also, as de Avillez & Breitschwerdt (2005)
have discussed on the basis of simulations of a magnetized ISM, su-
perbubble shells can stabilize against RT instability in the presence
of magnetic fields. In this regard, the clumping of the shell from
thermal instability at an early phase of evolution of the superbubble
can be important.

We have studied the evolution of superbubbles in stratified discs
analytically and with simulations. Our results can be summarized
as follows.

(i) Our analytic calculations show that radiation losses are im-
portant for superbubble dynamics. Radiation loss is more impor-
tant for superbubbles with continuous energy injection than a
 supernova remnant of similar total energy. We estimate almost
70 per cent of the total energy being radiated away. We have fur-
ther checked our analytical results with numerical simulations. We
found that analytic results match the simulations well, differing at
most by a factor of the order of unity for the case of large am-
bient density. The results obtained by the analytic means there-
fore provide a useful benchmark to compare with realistic simu-
lations. Also, for discs with large gas density, with
n0 ≥ 1 cm−3, superbubble breakouts are not possible for surface density of OB
stars ≤100(20/100 pc) kpc−2, or an equivalent energy injection sur-
face density of ≤(2–5) × 10−3 erg cm−2 s−1.

(ii) Superbubbles that emerge from the disc with Mach number
of the order of 2–3 require an energy injection rate of
∼10−4 erg cm−2 s−1, corresponding to explosions triggered by the
largest OB associations with 104 M⊙. This energy injection scale
connects to disc galaxies with synchrotron emitting gas in the
extra-planar regions.

(iii) Vigorous superbubbles that break out of the disc with suffi-
ciently large Mach number (≥10), correspond to an energy injection
rate of ∼10−3 erg cm−2 s−1, or equivalently, an SFR surface density of
∼0.1 M⊙ yr−1 kpc−2. These superbubbles require more than one
OB associations to produce and sustain their dynamics, and this
energy injection scale corresponds to (a) the existence of multiphase
gas in the halo of disc galaxies and (b) the Heckman threshold for
the onset of superwinds.

(iv) Superbubbles do not accelerate until reaching a vertical dis-
tance of a few scaleheights (of the order of ~2), which implies that
RT instability helps to fragment the shells not at a distance of a
scaleheight but at a much larger height. Also, we find that for typi-
cal disc parameters, thermal instability acts on the shell at the early

stages of superbubble evolution, and forms clumps and fragments
in the shell, much before the shell is acted upon by RT instability.
Radiative cooling therefore manifests in seeding thermal instability,
which has important implications for the clumping of superbubble
shell and producing channels of leakage for ultraviolet radiation
into the halo.

ACKNOWLEDGEMENTS

We thank Sergiy Silich for helpful comments on a draft of the paper.
We also thank an anonymous referee for the useful comments.
This work is partly supported by an Indo-Russian project (RFBR grant
08–02–91321, DST-India grant INT-RFBR-P121).

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