THE ORIGIN AND KINEMATICS OF COLD GAS IN GALACTIC WINDS: INSIGHT FROM NUMERICAL SIMULATIONS

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Received 2008 March 14; accepted 2009 April 6; published 2009 May 22

ABSTRACT

We study the origin of Na i-absorbing gas in ultraluminous infrared galaxies motivated by the recent observations by Martin of extremely superthermal linewidths in this cool gas. We model the effects of repeated supernova explosions driving supershells in the central regions of molecular disks with \( M_d = 10^{10} M_\odot \), using cylindrically symmetric gas dynamical simulations run with ZEUS-3D. The shocked swept-up shells quickly cool and fragment by Rayleigh–Taylor (R–T) instability as they accelerate out of the dense, stratified disks. The numerical resolution of the cooling and compression at the shock fronts determines the peak shell density, and so the speed of R–T fragmentation. We identify cooled shells and shell fragments as Na i-absorbing gas and study its kinematics along various sightlines across the grid. We find that simulations with a numerical resolution of \( \lesssim 0.2 \) pc produce multiple R–T fragmented shells in a given line of sight that appear to explain the observed kinematics. We suggest that the observed wide Na i absorption lines, \( \langle v \rangle = 520 \pm 120 \) km s\(^{-1} \), are produced by these multiple fragmented shells traveling at different velocities. We also suggest that some shell fragments can be accelerated above the observed average terminal velocity of \( 750 \) km s\(^{-1} \) by the same energy-driven wind with an instantaneous starburst of \( \sim 10^9 M_\odot \). The mass carried by these fragments is only a small fraction of the total shell mass, while the bulk of mass is traveling with velocities consistent with the observed average shell velocity \( 330 \pm 100 \) km s\(^{-1} \). Our results show that an energy-driven bubble causing R–T instabilities can explain the kinematics of cool gas seen in the Na i observations without invoking additional physics relying primarily on momentum conservation, such as entrainment of gas by Kelvin–Helmholtz instabilities, ram pressure driving of cold clouds by a hot wind, or radiation pressure acting on dust.

Key words: galaxies: starburst – hydrodynamics – ISM: bubbles – ISM: jets and outflows – ISM: kinematics and dynamics – supernovae: general

1. INTRODUCTION

Nearly all starburst galaxies, regardless of mass, appear to drive large-scale gaseous outflows, or galactic winds (Heckman et al. 1990; Martin 1999). Measurements demonstrate that these metal-enriched winds transport interstellar gas and supernova ejecta into galactic halos (Martin et al. 2002). These winds are thought to influence the thermal and chemical evolution of the intergalactic medium and hence the formation of galaxies as well as their evolution.

From radio to X-ray frequencies, observations of starburst galaxies reveal outflowing gas over a very broad temperature range (Martin et al. 2002). However, all observed emission is relatively near the galaxy, within a projected separation of about 10 kpc, due to the radial density gradient of the wind and density-squared dependence of emission processes. Absorption-line measurements are more sensitive to low-density, cold gas. The number of detections of blue-shifted (i.e., outflowing) interstellar absorption lines in starburst galaxy spectra has grown by a large factor in recent years (Heckman et al. 2000; Rupke et al. 2002; Schwartz & Martin 2004; Martin 2005). The shortcoming of absorption line measurements is that they do not uniquely determine the distance between the galaxy and the absorbing material.

Numerical simulations of galactic winds can provide needed insight into where the absorption originates. Using simulations to interpret observations, and observations to constrain simulations, is probably the only way to really understand these complex outflows dynamically. Modeling the early evolution of a galactic wind as it blows out of its disk requires a numerical, rather than analytic, approach due to the importance of nonlinear hydrodynamic and thermal instabilities.

Supershells evolve with roughly spherical geometry until they grow to scales of the disk gas scale height (Tomisaka & Ikeuchi 1986, 1988; Mac Low & McCray 1988; Tenorio-Tagle & Bodenheimer 1988; de Young & Heckman 1994). The acceleration of the shell into the galactic halo causes it to fragment via Rayleigh–Taylor (R–T) instabilities (Mac Low et al. 1989). The hot, low-density bubble interior radiates inefficiently. The wind can sweep up new shells of ambient gas, which in turn fragment by R–T instability, leaving a broad region containing fragments of fast-moving cool gas.

The swept-up shell is driven by the thermal pressure of the interior \( P = \rho c_s^2 \), where \( c_s \) is the interior sound speed. After blowout, the hot gas expands freely through the fragmented shell, producing a supersonic, energy-driven wind with velocity \( v_w \). Although entrained shell fragments can still be accelerated by the ram pressure of the wind \( P_{\text{ram}} = \rho v_w^2 \), this appears to be a minor contribution to their total kinetic energy. This can be seen by comparing the velocity of a bubble expanding into a uniform medium at a radius of one scale height to the final shell fragment velocities, as reported, for example, by Mac Low et al. (1989). These are the same, to within a factor of 2.

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Properties of the cool gas in starburst winds have been estimated from observations of interstellar Na i lines in starburst galaxy spectra. Estimates of the total mass of cold gas in these outflows have large uncertainties at present due to line saturation and dust depletion. Nonetheless, it has been emphasized that the momentum of the cool flows appear to be somewhat less than the amount available from either supernova ejecta or the radiation field, at least for the most luminous starbursts (Rupke et al. 2005; Martin 2006). The same approximations, however, also yield kinetic energies for the cool outflow that are only a few percent (up to a few tens of percent) of the supernova energy. The same flows could also be driven by energy-conserving bubbles, with only a small fraction of the total energy in the bubble going to accelerating the cold gas.

The standard scenario used for interpreting starburst wind absorption is based on the simulation shown in Figure 11 of Heckman et al. (2000), which suggests that dense clouds are advected into the wind at the interface between the low-latitude disk and the wind, by Kelvin–Helmholtz instability. With a grid resolution of 4.9 pc, however, the clumps of dense gas are not fully resolved in that simulation, leaving artificially large clumps that completely stop fragmenting below ~6 zone size. We will show below that the R–T instability is suppressed in secondary shells at that resolution, as well, substantially changing the distribution of cool gas.

Recently, Cooper et al. (2008) performed three-dimensional simulations of starburst blowout through a galactic disk with a fractal density distribution. They injected energy at a rate proportional to local density, rather than identifying supernova sites and following the explicit evolution of their remnants. This leads to higher than physical radiative losses, so their results represent a lower limit to the effects of a starburst. They found that Hα-emitting gas comes from the gas dynamical stripping and fragmentation of existing interstellar clouds. This gas can reasonably also be identified as a potential source of absorption, although they did not address the question explicitly or extend the simulation to times long enough to directly model the absorption.

We instead identify the location of the absorbing gas in fragmenting shells of swept-up interstellar gas, using high-resolution two-dimensional simulations with resolution as small as 0.1 pc. Although in our models we compute only up to the time of blowout because of the small region covered by our computational grid, we use the ballistic approximation (Zahnle & Mac Low 1995; Fujita et al. 2004) to show that gravitational deceleration does not act strongly on the shells and fragments during the starburst duration (10–40 Myr) if their velocities exceed 50–200 km s\(^{-1}\) at blowout. This analysis is based, however, on an assumption that the bulk of their mass remains unablated by the wind blowing past them. Understanding the full history of shell fragments and clumps will require substantial further work.

We address the origin and kinematics of the cold wind as measured in the Na i λλ5890, 96 absorption lines. The observations pose three major questions. First, why do the absorption line widths tend to greatly exceed the thermal velocity dispersion of warm neutral gas? The average full width at half maximum (FWHM) of the dynamic component is 320±120 km s\(^{-1}\) in ultraluminous infrared galaxies (ULIRGs), while the line widths range from 150 to 600 km s\(^{-1}\) in luminous infrared galaxies. Second, why do the terminal velocities of the cold gas approach the escape velocities from the starburst galaxies (Martin 2005)? Third, what do the maximum and mean velocities measured in the line profiles really represent physically?

We use our models to pursue five investigations. First, we investigate how the absorption properties change with viewing angle. We specifically test whether multiple R–T fragmented shells along a line of sight can reproduce the broad line width seen in Na i absorption lines. Second, we vary the numerical resolution to demonstrate how increased resolution of radiative cooling behind the shocks, and so of shell fragmentation, affect the results.

Third, we make a more general parameter study addressing variations in the properties of the outflowing cold gas with starburst luminosity, the size of the starburst region, and gas surface density. Fourth, we can obtain insight into the complicated dynamics of multiphase outflows, particularly their dependence on the mass loading of the wind. We investigate mass-loading rates between ~1.7 and 120 M\(_\odot\) yr\(^{-1}\) and vary the mechanical luminosity of the starburst between 10\(^{41}\)–10\(^{43}\) erg s\(^{-1}\) to see what velocities are reached by the swept-up shells and their fragments. The observed X-ray temperatures vary little with starburst luminosity T ~ 10\(^7\) K, so the terminal wind velocities should vary little with luminosity (above some critical value required for blowout).

Finally, Heckman et al. (2000) argued that both Na i-absorbing and Hα-emitting gas cannot originate in the swept-up shells because of the lack of strong correlation between the widths of Na i absorption lines and Hα emission lines. For example, the outflow sources with very broad (400–600 km s\(^{-1}\)) Na i absorption lines have Hα emission-line widths ranging from 145 to 1500 km s\(^{-1}\). Although a full nebular emission calculation is well beyond the scope of this paper, we do discuss where the ionization front might reside for various ionizing photon luminosities. We study the kinematics of the low-ionization Na i-absorbing gas and photoionized Hα-emitting gas by separating them crudely, using the photoionization code of Abel et al. (1999).

The acceleration of shell fragments is sensitive to how well shell fragmentation is resolved. Applying adaptive mesh refinement (AMR) techniques to this problem can maintain high resolution in the shocked shells and clouds. This paper is the first step toward such an improved simulation. We compare the blowout problem run on a fixed grid to a similar problem run with an adaptive grid, focusing on the comparison to measured properties of cold gas in galactic winds.

In this paper, we describe our disk and star formation models in Section 2 and our numerical method in Section 3. We give the results of our parameter studies in Section 4 and discuss comparisons with observations in Section 5, followed by conclusions in Section 6. In the Appendix, we show the results of test simulations of blowout in a dwarf galaxy by ZEUS-3D (Stone & Norman 1992a; Clarke 1994) and SAGE (SAIC’s Adaptive Grid, Eulerian hydrocode; Kerbyson et al. 2001; Gittins et al. 2008).

2. DISK AND STAR FORMATION MODELS

The parameters of our starburst model are based on the properties of ULIRGs to facilitate comparison with Martin (2005, 2006). We use hydrodynamic simulations to model the effects of multiple supernova explosions in the central 200 pc × 100 pc region of the molecular disk of a ULIRG. Our model is an extension of the blowout model in dwarf galaxies described by
Mac Low & Ferrara (1999) and Fujita et al. (2003). Our fiducial numerical resolution is 0.2 pc, sufficient to resolve cooling behind the shocks, and so the fragmentation of the swept-up shells by R–T instability as well as possible within a reasonable computational time. To study the effects of numerical resolution, we use models with resolution ranging from 0.1 to 0.8 pc.

2.1. Disk

ULIRGs are starburst galaxies with infrared luminosity $>10^{12} L_\odot$, and are usually found in major mergers and interacting galaxies (Sanders et al. 1988). They are believed to go through starburst phases twice, when the gas in a galaxy with a prograde orbital geometry is tidally disturbed during the first encounter with another galaxy and when both galaxies meet again and finally merge (Mihos & Hernquist 1996; Murphy et al. 2001; Li et al. 2004). We choose to model a molecular gas-rich spiral galaxy on its first encounter with another galaxy of similar mass.

We set up a molecular disk with $M_g = 10^{10} M_\odot$ in a dark matter halo with $M_{\text{halo}} = 5 \times 10^{12} M_\odot$. CO observations of ULIRGs at both first and second passages show the presence of molecular disks with $M_g = (0.4–1.5) \times 10^{10} M_\odot$ (Sanders et al. 1988; Solomon et al. 1997), which is in the range found for gas-rich spiral galaxies. However, the emission originates in regions a few hundred parsecs in radius, yielding surface densities of $\sim (0.5–1) \times 10^2 M_\odot$ pc$^{-2}$, within which the molecular mass dominates the dynamical mass (Sanders et al. 1988; Solomon et al. 1997). At these high surface densities, molecular hydrogen will dominate (Blitz & Rosolowsky 2006), as it can form within a few million years in turbulent regions with densities over 100 cm$^{-3}$ (Glover & Mac Low 2007). The density of H$_2$ traced by CO emission is $\sim 500$ cm$^{-3}$, comparable to the envelope of giant molecular clouds, while a region of much higher density in ULIRGs is traced by HCN emission, $\sim 10^3$ cm$^{-3}$, comparable to star-forming cloud cores (Solomon et al. 1992).

We assume that the entire interstellar medium (ISM) is a scaled-up version of a normal galactic disk with the ambient densities a factor of $\sim 100$ higher, making even the intercloud medium a molecular region. Thus we assume that the surface density distribution of the molecular disk is exponential, with $\Sigma(R) = \Sigma_0 \exp(-R/R_d)$ where $R_d$ is a scale radius (see also $\Sigma(R)$ of Arp 220 by Scoville et al. 1997).

We choose to model a disk with a central surface density of $\Sigma_0 = 10^4 M_\odot$ pc$^{-2}$, with a disk scale radius $R_d = 0.7$ kpc. The disk is in hydrostatic equilibrium with a Navarro et al. (1997; hereafter NFW) halo potential, and a disk potential based on the thin disk approximation (Toomre 1963), since $M_{\text{dyn}} \approx M_g$. The NFW potential is

$$\Phi(x) = \frac{G M_{\text{halo}} \ln(1 + cx)/x}{R_v F(c)},$$

where we set the virial radius $R_v = 326$ kpc, $x = r/R_v$, $c$ is a halo concentration factor, set to $c = 5$, appropriate for a large halo (Jimenez et al. 2003), and $F(c) = \ln(1 + c) - c/(1 + c)$.

The velocity dispersion of the molecular gas is observed to be 90 km s$^{-1}$ in Arp 220, which appears to be at the end of the merging process (Scoville et al. 1997). Such a high velocity dispersion yields a scale height of 15 pc in its disk (Scoville et al. 1997). Molecular clouds with such high dispersion will get destroyed by colliding with other clouds, so the cooling time behind the shocks must be shorter than the destruction time interval. We assume the gas is supported by turbulence with a similarly high velocity dispersion $c_v = 55$ km s$^{-1}$. It is lower than 90 km s$^{-1}$ because gravity from our disk gas and halo cannot confine the gas with a higher velocity dispersion. As a comparison, we also model a disk with a higher surface density of $\Sigma_0 = 5 \times 10^4 M_\odot$ pc$^{-2}$ with $R_d = 0.17$ kpc and $c_v = 90$ km s$^{-1}$, based on the central surface density observed in Arp 220. This is the highest $\Sigma_0$ of all observed ULIRGs.

Figure 1 shows the vertical density distributions of both disks. The exponential scale heights are rather small, 7 and 2 pc, but the gas within them is very dense, $\sim 500$ and $10^3$ cm$^{-3}$, respectively. The gas density is still $\sim 500$ cm$^{-3}$ at $Z \geq 4R_d$. At higher altitudes, where the gas is less dense, the gas is physically atomic or even ionized. We do not take that state change into account in our model, though. When the number density drops to $n = 10^{-2}$ cm$^{-3}$ we set the gas density in the halo constant as it is no longer dynamically important on the length and timescales treated in our model.

2.2. Star Formation

We assume a single starburst that occurs at the center of the disk, and that all the kinetic energy of the starburst supernovae is released in a central wind of constant mechanical luminosity. In reality, the discrete energy inputs from supernovae generate blastwaves that become subsonic in the hot interior of the bubble first produced by stellar winds, and hence can be treated as a continuous mechanical luminosity in the study of bubble dynamics (Mac Low & McCray 1988). These assumptions mean that a single superbubble forms, evolving to produce a bipolar outflow of gas.

Figure 2 shows the evolution of mechanical luminosity $L_{\text{mech}}$ as a function of time for an instantaneous starburst with $10^7 M_\odot$ of gas turning into stars and for continuous starbursts with star formation rates of 100 and 500 $M_\odot$ yr$^{-1}$, based on the Starburst 99 model (Leitherer et al. 1999). The Starburst 99 model uses a power-law initial mass function with exponent $\alpha = 2.35$ between low-mass and high-mass cutoff masses of $M_{\text{low}} = 1 M_\odot$ and $M_{\text{up}} = 100 M_\odot$ with solar metallicity. Star formation rates in ULIRGs are estimated to be $\gtrsim 100 M_\odot$ yr$^{-1}$ based on far-infrared luminosities and the assumption of continuous star formation (see Table 1 of Martin 2005; note that...
the star formation rates given there correspond to a low mass
cutoff of 0.1 $M_\odot$ and must be divided by a factor of 2.55 before
corresponding to an instantaneous starburst with $M_* = 10^6 M_\odot$.
For a continuous starburst, $L_{\text{mech}}$ increases until the death rate
of massive stars catches up to their birth rate, after about 40 Myr.
The power rises particularly rapidly over the first few Myr, the period
modeled by our simulation. Figure 2(a) shows the evolution for a
continuous star formation with 500 $M_\odot$ yr$^{-1}$ and 100 $M_\odot$ yr$^{-1}$.
For our ULIRG models, we use constant mechanical luminosity winds
with values $L_{\text{mech}} = 10^{43}$, $10^{42}$, and $10^{41}$ erg s$^{-1}$. The highest
of these corresponds to the mechanical luminosity expected from
stellar winds during the first 2 Myr of an instantaneous starburst
with $M_* = 10^6 M_\odot$. The subsequent supernovae will result in a
far higher mechanical luminosity, but that is likely to be vented
out of the galactic disk through the hole opened by the initial
blowout. $L_{\text{mech}} = 10^{43}$ erg s$^{-1}$ also corresponds to a model
with a constant star formation rate (SFR) of 15.9 $M_\odot$ yr$^{-1}$. This
correspondence assumes that the birth and death rates of massive
stars are in equilibrium, which is achieved about 40 Myr after
the burst begins. In this scenario, the burst would be ongoing
through its initial stages before our simulation starts, but gas
displaced by the initial feedback had been replaced by inflows,
and prior feedback energy was largely radiated. Our simulation
is not such a good representation of this scenario.

We note it is an oversimplification to assume a single starburst
at the disk center for modeling a galactic outflow. Bipolar
outflows are seen in some starburst galaxies such as M82
(Strickland & Stevens 2000; Strickland et al. 2004) for example,
although some ULIRG winds appear to require starburst regions
extended to $\gtrsim 1$ kpc to launch the cool outflow (Martin 2006).
Three-dimensional, hydrodynamic simulations of dwarf
starbursts have shown that extended, multiple energy sources,
as well as a single central energy source, form a bipolar outflow
(Fragile et al. 2004). These simulations also demonstrated that
the main effect of multiple sources was to reduce the fraction
of metals and energy ejected from the galaxy from almost unity
to around 50%. Our assumption thus represents a reasonably
strong lower bound to the amount of kinetic energy that will be
deposited in the observed cold gas. We therefore start with this
assumption and do not expect the results to differ much from
those expected with more extended star-forming regions.

In addition, we neglect the effects of UV radiation on
molecular hydrogen in the disk. The UV radiation from massive
stars may photodissociate some of molecular hydrogen outside
star-forming cores to atomic hydrogen. However, the assumed
turbulent pressure with $c_s = 55$ km s$^{-1}$ is 14 times greater
than the increased thermal pressure by photodissociation. We
thus safely neglect the effects of UV radiation on the disk gas
structure.

We define the model with $\Sigma_0 = 10^4 M_\odot$ pc$^{-2}$ and $L_{\text{mech}} =
10^{43}$ erg s$^{-1}$ as our fiducial model (U1/X1). We list the
parameters for all the other runs in Table 1.

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**Figure 2.** Mechanical luminosities (left panel) and ionizing photon luminosity (right panel) as a function of time for three starburst scenarios: instantaneous starburst
with $M_* = 10^6 M_\odot$ (solid line); continuous star formation with 500 $M_\odot$ yr$^{-1}$ (dashed line); and continuous star formation with 100 $M_\odot$ yr$^{-1}$ (dash-dotted line).
Population synthesis models are from Starburst 99 (Leitherer et al. 1999).
Figure 3. (a) Temperature of a hot bubble in a uniform medium driven by a starburst with $L_{\text{mech}}$ and the rate of supernova ejecta $M_{SN} \propto L_{\text{mech}}$ predicted by the Starburst 99 model as a function of time, with the given mass-loading factors $\xi$, under the assumption of an instantaneous (solid line) or continuous (dashed line) starburst. The amount of mass in the hot wind is $\xi M_{SN}$. (b) The expected terminal velocity of the wind driven by such a bubble after its blowout from a stratified disk. Note $T_{\text{wind}}$ and $v_{\text{wind}}$ are proportional to $L_{\text{mech}}/M_{SN}$, thus independent of SFR assumed.

### 2.3. Bubble Dynamics

In the molecular disk of a ULIRG with a very small scale height, a bubble with $L_{\text{mech}} \geq 10^{41}$ erg s$^{-1}$ quickly blows out of the disk at $t \ll 1$ Myr. We can only simulate the evolution of the bubble up to $t \sim 0.3$–1 Myr before it leaves our grid, which is rather small because of the cost of high resolution. We argue in Section 5 that cooled, swept-up shells, which we identify as Na i-absorbing gas, can acquire maximum velocities primarily determined by when they blow out, accelerate, and fragment. The question remains whether these shell fragments and clumps survive as the hot interior wind streams through them.

Each individual shell fragment after blowout remains unstable to smaller scale R–T instabilities while being further accelerated by the wind, requiring extremely high resolution to fully resolve. Resolution is not as critical for previously published numerical studies on starburst winds that explored feedback parameters. For example, reasonable assumptions about cooling losses indicate that moderate luminosity starbursts do not remove a significant fraction of the galaxy’s gas from the halos (Mac Low & Ferrara 1999), and the mass loss rates would only decrease further with more cooling. A large fraction of the heavy elements do escape from the halos (Mac Low & Ferrara 1999; Fujita et al. 2004), and this result does not depend on resolution below ~ a few tens of parsecs, so long as at least a few R–T modes are excited in the swept-up shells to let the metal-enriched gas escape. Therefore, a major question that must be answered to understand observations of cold gas at high velocity is whether clumps of dense gas survive in the lower density wind (see Klein et al. 1994; Marcolini et al. 2005), where they will fragment because of hydrodynamic instabilities, such as R–T (strictly speaking, Richtmyer–Meshkov; see Richtmyer 1960; Meshkov 1969) and Kelvin–Helmholtz.

Weaver et al. (1977) argued that the density in the hot interior of bubbles in uniform gas is dominated by conductive evaporation from the dense shell. However, our model does not explicitly include thermal conduction, or material ablated off of high-density molecular clouds associated with the central starburst. Instead, we add additional mass to our central luminosity source to account for this process, multiplying the mass input rate by a mass-loading factor $\xi$. The amount of thermally evaporated mass in a bubble expanding into a uniform medium is proportional to $L_{\text{mech}}^{2/35} \rho^{-2/35}$. The internal temperature is

$$T_b(t) = (\gamma - 1) \frac{\mu}{k_B} \frac{5}{11} \int_0^t L_{\text{mech}}(t') dt' \frac{M_{SN}}{\xi M_{SN}(t')} \varepsilon(t'),$$

where $L_{\text{mech}}(t)$ and the mass of supernova ejecta $M_{SN}(t)$ as a function of time are taken from Starburst 99 model, the adiabatic index $\gamma = 5/3$, mean mass per particle $\mu = 14/22 m_H$, and $k_B$ is the Boltzmann factor. Weaver et al. (1977) showed that $5/11$ of the total input mechanical energy goes into the hot, pressurized region bound by the inner and outer shock fronts (see their Equation (14)). This was applied to superbubbles by McCray & Kafatos (1987). We show $T_b(t)$ in Figure 3(a).

We can estimate the terminal velocity of the wind driven by such a bubble by equating the bubble interior energy with the kinetic energy of the mass-loaded wind, so that

$$v_{\text{wind}}^2(t) = 2 \left[ \frac{5}{11} \int_0^t L_{\text{mech}}(t') dt' \right] / \left[ \int_0^t \xi M_{SN}(t') dt' \right],$$

which we plot in Figure 3(b). Since the mechanical luminosity $L_{\text{mech}}(t)$ and the mass of supernova ejecta $M_{SN}(t)$ are both linearly proportional to the amount of gas converted to stars, $T_b$ and $v_{\text{wind}}$ are the same for all strengths of starburst with the same mass-loading rate $\xi$. We take a fiducial $\xi$ value of about 8 corresponding to a mass-loading rate $M_{\text{in}} = 17(L_{\text{mech}} / 10^{43}) M_\odot \text{yr}^{-1}$, but run test simulations with $\xi$ varying between ~1 and 15 to explore the effects of mass loading on the shell kinematics. Comparisons to observations suggest $\xi \approx 10$ (Suchkov et al. 1996—see M82 comparison; Martin et al. 2002). Throughout this paper, we designate as the wind the hot interior gas freely streaming outward after blowout of the shell.

### 3. NUMERICAL METHODS

We follow the numerical methods used by Mac Low & Ferrara (1999) and Fujita et al. (2003) to model a starburst...
in a galactic disk. We briefly summarize our methods below, but refer to the papers above for more details. We compute the evolution of a starburst-driven blastwave as it blows out of the molecular disk of a ULIRG with ZEUS-3D,8 a Eulerian, finite-difference, astrophysical gas dynamics code (Stone & Norman 1992a; Clarke 1994), that uses second-order van Leer (1977) advection, and a quadratic artificial viscosity to resolve shock fronts. We use the loop-level parallelized version ZEUS-3D, in its two-dimensional form. Runs were done on Silicon Graphics Origin 2000 machines using eight processors, and typically took ~1–12 days.

We assume azimuthal symmetry around the rotational axis of the galaxy. Our fiducial grids are 1000 × 500 zones with a resolution of 0.2 pc, comparable to the size of star-forming cloud cores. We also ran the same simulations with resolution of 0.1, 0.4, and 0.8 pc to examine the sensitivity of post-shock shell density and thus R–T instability growth to resolution. We use reflecting boundary conditions along the symmetry axis and along the galaxy midplane and outfall boundary conditions on the other two axes.

The assumption of azimuthal symmetry limits R–T instabilities to growing as rings, reducing the number of clumps below what would actually be expected from three-dimensional spikes. Mac Low et al. (1989) compared models with azimuthal to slab symmetry to demonstrate that the fixing of a central axis of symmetry did not markedly change the behavior. Three-dimensional models of isolated superbubbles have only been performed for the magnetized case starting with the work by Tomisaka (1998). Recently Stil et al. (2008) have performed three-dimensional studies of unmagnetized superbubbles as calibration for a study of magnetized superbubbles, but they did not extend their hydrodynamical models into the R–T unstable regime. Comparison of two-dimensional to three-dimensional models of shocked clouds by Stone & Norman (1992a) and Xu & Stone (1994) found little difference in their dynamical evolution aside from the breakup of post-shock vortex rings in three dimensions. Young et al. (2001) and Cabot (2006) compared high resolution two-dimensional and three-dimensional models of planar, incompressible R–T instability. Cabot (2006) cautions that two-dimensional models produce larger, less well mixed structures at late times because of the inverse energy cascade that occurs in two-dimensional flows.

To drive a constant luminosity wind, we add mass and energy to a source region with a radius of 10 pc (50 zones). Our fiducial mass input rate is \( \dot{M}_{\text{in}} = 17(L_{\text{mech}}/10^{43}) M_\odot \text{yr}^{-1} \) which corresponds to a mass-loading factor \( \xi \approx 8 \). For simulations with different resolutions, we keep the number of zones covering the spherical edge of the source regions the same by maintaining its radius as a constant number of zones rather than a constant physical size. This is important because aliasing at the edge creates density perturbations that can be amplified by hydrodynamic instabilities, such as the R–T instability in the swept-up shell. We use ratioed grids for lower-resolution runs and we decrease the size of source region to a radius of 5 pc (50 zones) for a higher-resolution run. We directly show the effects of this initial noise on the development of R–T instability by running a simulation with a resolution of 0.2 pc, but with a source region with a radius of 5 pc (25 zones).

As in Mac Low & Ferrara (1999), we use a cooling curve by MacDonald & Bailey (1981) for solar metallicty with a temperature floor of either \( T_{\text{floor}} = 10^2 \) K, the temperature to which the metals can cool the gas, or \( 10^4 \) K, the temperature maintained by photoionization heating. The shocked gas in swept-up shells cools efficiently to the temperature floor set in the cooling curve, since our molecular disk is very dense. We show in Section 4.2.2 that even our highest resolution runs do not yet fully resolve the dense shells even for \( T_{\text{floor}} = 10^4 \) K, so the influence of the cooling floor is not evident in our work. We also include an empirical heating function tuned to balance the cooling in the background atmosphere. This is linearly proportional to density, so that it is overwhelmed by cooling in compressed gas which is proportional to the square of the density (Mac Low et al. 1989). This is to prevent the background atmosphere from spontaneously cooling. We use a tracer field (Yabe & Xiao 1993) to turn off radiative cooling in the hot bubble interior, in order to prevent mass numerically diffused off the dense shell from spuriously cooling the interior. The cooling time of the interior is much longer than the dynamical time of our bubble, so interior cooling is physically unimportant to the bubble dynamics (Mac Low & McCray 1988). These adiabatic bubble interiors form energy-driven winds after blowout.

4. PARAMETER STUDIES

We now describe the results of parameter studies of both physical and numerical variables. We begin by considering physical variables, including mechanical luminosity, mass loading of the wind, and disk surface density. We then discuss numerical variables, focusing on how numerical resolution and the cooling cutoff temperature affect shell density and also examining the effect of changing the size of the source region.

4.1. Physical Parameters

4.1.1. Mechanical Luminosities

Our fiducial model (X1) has mechanical luminosity \( L_{\text{mech}} = 10^{43} \) erg s\(^{-1}\). This model corresponds to the first 2 Myr of a starburst in which 10\(^9\) \( M_\odot \) of gas turns into stars instantaneously. We compare this to models with lower mechanical luminosities \( L_{\text{mech}} = 10^{42} \) and \( 10^{44} \) erg s\(^{-1}\) (models X2 and X3, respectively) in the same molecular disk. These mechanical luminosities correspond to instantaneous bursts of 10\(^8\) and 10\(^7\) \( M_\odot \). Figure 4 shows the density distribution of our fiducial model in its right panel. This may be compared to Figure 5, which shows the density distributions of the two models with lower \( L_{\text{mech}} \) at \( t \approx 0.49 \) Myr and 0.85 Myr, respectively. The swept-up shells fragment due to R–T instability into multiple clumps and shells.

Secondary Kelvin–Helmholtz instabilities ablate the sides of these fragments as the hot gas streams through them. Look at the clumps, for instance, at (R,Z) of (40,65), (30,110), and (15–25,120) pc in X1. Note also that the swept-up shells in the horizontal direction are also R–T unstable, because our disk gas is stratified in the radial direction, too, due to its exponential surface density profile.

In fact, the degree of fragmentation is larger in X2 and more so in X3 because \( \dot{M}_{\text{in}} \propto L_{\text{mech}} \) and so the density of interior gas is lower. In particular, most shell fragments in X3 are already falling back to the disk. A mechanical power of \( 10^{46} \) erg s\(^{-1}\) is just too small in such a dense environment to accelerate the bulk of the shells to the disk’s escape velocity.

4.1.2. Surface Density

The top right panel of Figure 5 shows the density distribution of our model with \( M_d = 10^{10} \) \( M_\odot \), but with a higher surface...
density, $\Sigma_0 = 5 \times 10^4 \, M_\odot \, \text{pc}^{-2}$ at $t = 0.22$ Myr (V1). With the same mechanical luminosity, $L_{\text{mech}} = 10^{43}$ erg s$^{-1}$, the bubble blows out earlier at $t \approx 0.15$ Myr, because the disk is more stratified with a smaller scale height (see Figure 1). Except the time of blowout, the degree of fragmentation and the shell kinematics are about the same in models with surface densities different by a factor of 5.

4.1.3. Mass Loading

Mass loading from thermal conduction and molecular clouds determines the density of the bubble interior and wind. Figure 6 shows the density distributions of our fiducial ULIRG model with $T_{\text{floor}} = 10^2$ K with mass-loading rates of 1.7, 17, 49, and 120 $M_\odot$ yr$^{-1}$ (models U1-A, U1, U1-B, and U1-C). These mass-loading rates correspond to bubble interior temperatures, $T_b = 1.7 \times 10^8, 2.2 \times 10^7, 7.5 \times 10^6,$ and $3.2 \times 10^6$ K and wind terminal velocities expected when all the thermal energy is converted to kinetic energy, $v_{\text{wind}} = 2700, 1000, 600,$ and 250 km s$^{-1}$, respectively.

Figure 6 shows the bubbles just before they leave the grid, at $t = 0.23, 0.27, 0.35,$ and 0.41 Myr, respectively. Since the input mechanical luminosity is the same in all these models the bubbles initially grow at about the same rate, driven by the thermal pressure of the hot interior gas. However, a bubble with a lower mass-loading rate and higher wind terminal velocity expands faster into the halo once the swept-up shells fragment and the hot gas blows out between the fragments. In addition, the blowout occurs earlier and in more places with a lower mass-loading rate because the density of the bubble interior gas is lower. A higher density contrast between the hot interior and the swept-up shells promotes shell fragmentation by R–T instability, the growth of which is proportional to the density contrast. In the least dense model U1-A, all the swept-up shells quickly fragment into fingers and filaments and the dense clumps at their edges are subject to secondary Kelvin–Helmholtz instabilities.

As the hot bubble interior becomes transonic after blowout, it accelerates the shells and shell fragments by ram pressure rather than thermal pressure. This low-density wind thus can acceler-
They are shown at $t = 0.22, 0.27, 0.35, \text{ and } 0.41 \text{ Myr, respectively. Different mass-loading rates correspond to different wind temperatures and terminal wind velocities. Note that the growth of hydrodynamic instabilities is suppressed as the mass-loading rate decreases and the terminal wind velocity increases.}

4.2. Numerical Parameters

4.2.1. Grid Resolution

Figure 7 shows the density distributions of our fiducial model with grid resolution varying from 0.1 to 0.8 pc. We chose $T_{\text{floor}} = 10^4 \text{ K}$ for this resolution study, because shell densities are not too far from what we expect analytically with this high minimum temperature. The growth of R–T instability is significantly enhanced in the highest resolution run, X1-0, and suppressed in the lower resolution runs, X1-2 and X1-4. In particular, all the outermost shells seen in X1 are further fragmented by R–T instability in X1-0 with the resolution increased by a factor of only 2. The positions of outer shock fronts in the horizontal direction agree very well among the four simulations, since the shells there are not subject to severe hydrodynamic instabilities.

Figure 8 shows the density profiles at the outer shock fronts in the vertical direction at $t = 0.06 \text{ Myr for } 0.1 \text{ pc resolution (triangles; model X1-0), and at } t = 0.05 \text{ Myr for } 0.2 \text{ pc (stars; X1), 0.4 pc (diamonds; X1-2), and 0.8 pc (crosses; X1-4) resolution. The shell is better resolved with a higher resolution, but still not fully resolved even at } 0.1 \text{ pc resolution. Note that the size of the highest resolution bubble is slightly smaller only because of the smaller source region.}

Figure 6. Density distributions of our model with $T_{\text{floor}} = 10^2 \text{ K}$ with different mass-loading rates: $M_\text{in} = 1.7 M_\odot \text{ yr}^{-1}$ (U1-A: top left), 17 $M_\odot \text{ yr}^{-1}$ (U1: top right), 49 $M_\odot \text{ yr}^{-1}$ (U1-B: bottom left), and 120 $M_\odot \text{ yr}^{-1}$ (U1-C: bottom right). They are shown at $t = 0.22, 0.27, 0.35, \text{ and } 0.41 \text{ Myr, respectively. Different mass-loading rates correspond to different wind temperatures and terminal wind velocities. Note that the growth of hydrodynamic instabilities is suppressed as the mass-loading rate decreases and the terminal wind velocity increases.}

Figure 7. Density distributions of our standard model with $T_{\text{floor}} = 10^4 \text{ K at } t = 0.27 \text{ Myr, with resolution of } 0.1 \text{ pc (X1-0: top left), our fiducial resolution of } 0.2 \text{ pc (X1: top right), and resolutions of } 0.4 \text{ pc (X1-2: bottom left) and } 0.8 \text{ pc (X1-4: bottom right). Note that the growth of R–T instability is suppressed as the resolution decreases. The black lines show the typical line of sight used for line profile analysis.}

Figure 8. Density profiles at the outer shock fronts in the vertical direction at $t = 0.06 \text{ Myr for } 0.1 \text{ pc resolution (triangles; model X1-0), and at } t = 0.05 \text{ Myr for } 0.2 \text{ pc (stars; X1), 0.4 pc (diamonds; X1-2), and 0.8 pc (crosses; X1-4) resolution. The shell is better resolved with a higher resolution, but still not fully resolved even at } 0.1 \text{ pc resolution. Note that the size of the highest resolution bubble is slightly smaller only because of the smaller source region.}

The observed X-ray temperature $T_X$ is $\sim 0.67 \text{ keV} = 7.7 \times 10^6 \text{ K}$ in all kinds of starburst galaxies from dwarfs to ULIRGs (Martin 1999; Heckman et al. 2001; Huo et al. 2004; Grimes et al. 2005). This corresponds to $\xi \approx 10$ in Equation (3). However, the X-ray emission is proportional to $n^2$ so it is biased toward high-density regions such as the interface between the hot interior gas and the shells and their fragments. At this interface, conductive evaporation and turbulent ablation raise the density. The recent observations of diffuse hard X-ray emission in starburst galaxies suggest the existence of a very hot $(\log T \gtrsim 7.5)$ metal-bearing gas (e.g., Strickland et al. 2004; Strickland & Heckman 2007). Then the bulk of the hot wind may still be very hot $\sim 10^8 \text{ K},$ the temperature which Strickland & Stevens (2000) modeled for M82 with $\xi = 1$. 

4.2. Numerical Parameters
of the source region in X1-0 to be half of that in X1 in order not to overproduce noise at the contact discontinuity. We will show below that this noise feeds R–T instability and must be maintained the same in order to study the effects of resolution on shell fragmentation alone. Thus the density profile of X1-0 at \( t = 0.06 \) Myr is not directly comparable to those of other runs at \( t = 0.05 \) Myr, but Figure 8 still demonstrates the trend in resolving the peak shell density as a function of resolution.

Figure 8 shows that the shell density is progressively better resolved as the resolution increases. However, the cooling times at the shock front are typically of order \( 10^2 \) yr (see Section 4.2.2), so a shock with cooling floor equal to the background temperature of \( 10^4 \) K may be treated as isothermal. At the time displayed, the Mach number of the outer shock is \( M = 19 \). The shell density we expect from an isothermal shock propagating into a background density \( \rho_{bg} \) is \( \rho_{sh} M^2 = 6.1 \times 10^{-10} \) g cm\(^{-3}\). Figure 8 shows that the shell density \( \rho_{shell} \) is still unresolved by a factor of \( \sim 4 \) even with our highest grid resolution of 0.1 pc.

The difference we see in the development of R–T instability develops because of two factors. First, the postshock density of the swept-up shells is not fully resolved. Increasing density contrast drives faster R–T growth. Second, the linear R–T instability develops because of two factors. First, the postshock density of shocked shells is approximately the same in both simulations. We will show below that this noise feeds R–T instability and must be not to overproduce noise at the contact discontinuity. We will show below that the resolution of 0.2 pc is still not sufficient to properly model hydrodynamic instabilities acting on shells and clumps, but is just sufficient for the purpose of demonstrating a wide range of velocities in shell fragments caused by R–T instability.

### 4.2.2. Cooling Temperature Floor

We show the density distributions of our fiducial model with different cooling curve cutoffs of \( T_{\text{floor}} = 10^2 \) K (U1) and \( 10^4 \) K (X1) at the time of blowout in Figure 4. Figure 4 shows that the temperature floor we choose for our cooling function appears to have a negligible influence on the evolution of bubbles in our models.

However, closer examination reveals that the density of shocked shells is approximately the same in both simulations despite the difference in temperature floor. For example, at \( t = 0.05 \) Myr before any fragmentation occurs, the shell density in both simulations is \( \sim 8 \times 10^{-20} \) g cm\(^{-3}\) in the vertical direction where the background disk density is \( \rho_{bg} = 2 \times 10^{-21} \) g cm\(^{-3}\). This is because the density peak in the simulations is limited by resolution in these models, not by the strength of cooling. The shock velocity in the vertical direction at \( t = 0.05 \) Myr is \( \sim 230 \) km s\(^{-1}\), or Mach number \( M = 15 \) in \( 10^4 \) K background gas. The immediate postshock temperature is then \( T = 7.6 \times 10^8 \) K. This shocked gas quickly cools to or below \( 10^4 \) K because the exponential cooling time (e.g., Mac Low & McCray 1988) is very short,

\[
\tau_{\text{cool}} = \frac{3kT}{4n_{bg} \Lambda} \approx 64 \text{ yr},
\]

with the mean mass per particle \( \mu = 14/22m_H \) for ionized gas, and \( \Lambda(T) = 4.1 \times 10^{-25} \) erg cm\(^3\) s\(^{-1}\) from the MacDonald & Bailey (1981) cooling curve.

The shell density expected from an isothermal shock will then be \( \rho_{shell} = \rho_{bg} M^2 \approx 5 \times 10^{-19} \) g cm\(^{-3}\) with Mach number \( M = 15 \) if \( T_{\text{floor}} = 10^4 \) K. If \( T_{\text{floor}} = 10^2 \) K, the shell density will reach values even greater than the isothermal value. However, since the shell density is far from being resolved even with 0.1 pc resolution in our simulations, the cooling floor has a negligible influence on our models.

### 4.2.3. Effects of Source Region on Shell Fragmentation

The top left panel of Figure 5 shows the density distribution of our fiducial model with \( T_{\text{floor}} = 10^4 \) K, but with a source region of a smaller radius of 5 pc or 25 zones (S1). The density structure of S1 looks very similar to that of X1, but the fragmentation by R–T instability is slightly suppressed. This is because a smaller number of cells is covering the edge of the spherical source region in which the density is imperfect. This imperfection creates a perturbation which gets amplified by hydrodynamic instabilities. A much smaller source region will make a bigger difference in the amount of fragmentation. It is important to note that we are anyway probably underestimating the degree of fragmentation since our bubble sweeps up a smooth, one-phase ISM instead of real ISM with density fluctuating strongly from the mean (see Cooper et al. 2008), and we have only a small, central source region instead of an extended distribution of supernovae (cf. Fragile et al. 2004).
5. COMPARISON TO OBSERVATIONS

In this section, we first use the ballistic approximation to justify making a comparison between the observations and our simulations in which the bubbles are evolved only up to $t \ll 1$ Myr. Then, we identify gas parcels likely to produce Na i absorption, and simulate observations of this gas along sightlines toward the galactic nucleus. The velocity spread, the mass-weighted velocity, and the maximum velocity are compared for different viewing angles. The velocity gradient across one of these winds is also studied for comparison to longslit spectra. We then compare our models to the observed Na i absorption profiles.

5.1. Ballistic Approximation

We use a limited grid size in order to maintain high resolution, so we simulate the evolution of bubbles only up to the time of blowout. We showed in the previous section that the bubbles blow out very early, at $t \ll 1$ Myr, because of the small scale height of our molecular disk. For comparison, ultraluminous starbursts have ages of up to 50 Myr (Murphy et al. 2001). To extrapolate our computational results to later times, we use the ballistic approximation (Zahnle & Mac Low 1995; Fujita et al. 2004) that after blowout, shell fragments travel on radial ballistic orbits in the gravitational potential of the galaxy, with no further accelerations by gas pressure gradients.

The strength of this approximation depends on the behavior of the wind, and conditions in the region into which the wind penetrates. This approximation was successfully used by Zahnle & Mac Low (1995) to follow the ejecta of a typical Shoemaker-Levy 9 fragment falling back to Jupiter’s atmosphere, and was found to give results quantitatively consistent with the observations of the size and shape of the infrared bright spots. That case does differ from the starburst case in having neither an ongoing wind nor a complex density distribution above the site of the explosion. In the starburst case, continuing supernova explosions at late times drive an ongoing wind (see Figure 2), while the mergers that drive most starbursts yield a complex geometry above the site of the blowout. However, even if the stronger winds expected from supernovas at late times in an instantaneous starburst do accelerate the fragments further, the ballistic approximation will still yield a lower limit to their velocities. Since the question of most interest is whether cold gas can be accelerated to such high velocities, a lower limit is already useful. The complex geometry of tidal tails may also not be of great concern, as models suggest that they only cover a small fraction of the solid angle visible from the nucleus, and so are unlikely to be dynamically dominant.

Under the ballistic approximation, the equation of motion for a shell fragment at a distance $r$ from the galactic nucleus is

$$v(r) = \left[ v_b^2 + 2(\Phi(r_b) - \Phi(r)) \right]^{1/2},$$

(5)

where $\Phi(r)$ is the total halo and disk potential, and $v(r_b) = v_b$ is the shell velocity at blowout at a position of $r_b$.

In Figure 9, we plot $v(r)$ for several initial velocities starting at a fixed initial position, $r_b = 200$ pc, just above our model disk. By setting $r_b = 200$ pc, we can ignore the disk potential which is negligible above the disk compared to the halo potential. Thus we set $\Phi = \Phi_{halo}$ in Equation (5) and solve the equation analytically using Equation (1). The upper limit to the radius reached by a shell fragment at time $t$ is given by the linear approximation $r \leq v_b t$, because the shell will only decelerate in the potential. We show actual radii for different initial velocities at $t = 10$ Myr by vertical ticks in the left panel of Figure 9.

Although Figure 9 is plotted as a function of radial velocity, we find similar results if total velocity is used instead, because shell fragments are traveling in nearly radial directions. This figure shows that the linear approximation is quite good for $v_b > 100$ km s$^{-1}$, so we use it to follow the cold gas for times of order 10 Myr.

For times $t \gg 100$ Myr, the linear approximation fails and most gas falls back to the center after reaching heights of a few hundred kpc. Only gas with $v_b \gtrsim 1000$ km s$^{-1}$ can completely escape the potential of the disk and halo (the halo alone has an escape velocity $v_{esc} = 800$ km s$^{-1}$ at radii less than 0.01$R_s$). We show below that very little gas actually escapes the halo, with $\lesssim 0.1\%$ of the total shell mass accelerated above $v_b \gtrsim 1000$ km s$^{-1}$ in our fiducial model. Our results confirm the result from previous simulations of smaller galaxies that the loss of ISM mass is inefficient (e.g., Mac Low & Ferrara 1999; D'Ercole & Brighenti 1999). However, significant mass is circulated over scales of 10 kpc, presenting a significant absorption cross section as suggested by Martin (2006).

5.2. Cool Gas

Only the coolest gas can be observed in Na i absorption. We chose a temperature cutoff of $T < 5 \times 10^4$ K to trace this gas because cooling remains numerically limited even with our highest resolution grid. Gas below this cutoff is so dense that the cooling time (Equation (4)) is very short, $\tau_{cool} < 0.01$ Myr, so it is physically expected to reach temperatures where Na i is present. As an example, the least dense shell fragment in Figure 10 has $\rho = 3.0 \times 10^{-24}$ g cm$^{-3}$ and $T = 3.3 \times 10^4$ K. Equation (4) gives an exponential cooling time $\tau_{cool} \approx 0.007$ Myr, using $A(T) \approx 4.2 \times 10^{-23}$ erg cm$^{-2}$ s$^{-1}$ from the MacDonald & Bailey (1981) cooling curve.

Figures 10 and 11 show the gas density, temperature, and radial velocity in models X1 and X1-0 along a line of sight through the central continuum source at an angle $\theta = 19^\circ$ from the vertical axis. Note that our temperature cutoff picks up only the densest and the coolest parts of the shells excluding numerically diffused interfaces between the shells and the hot gas.

We compare the characteristics of this temperature-selected gas to the measured properties of Na i absorption in ultraluminous starbursts. As argued above in Section 2.2, our assumption of a central point source of wind luminosity likely represents a lower limit to the amount of cold gas entrained. Fragile et al. (2004) found that, in galaxies where a central energy source
Figure 10. In the bottom left panel, the distribution of column density for the cold gas as a function of radial velocity is plotted. The other three panels plot radial profiles of radial velocity (top left), density (top right), and temperature (bottom right) along a line of sight through the center at an angle of 19° from the vertical axis in model X1 at the end of the simulation. The radial profiles all use the horizontal axis labeled on the bottom right. Regions of cold gas with $T < 5 \times 10^4$ K (which we take to be Na i-absorbing gas) are shown in diamonds.

Figure 11. Same as in Figure 10 for the same standard simulation, but with our highest resolution of 0.1 pc (X1-0). Again, the three radial profiles (top right, top left, bottom right) all use the radial axis given on the bottom right.

Figure 10 and Figure 11 show the distribution of column density, radial velocity, density, and temperature for cold gas in Galactic winds. The bottom left panel in Figure 10 highlights the distribution of column density as a function of radial velocity. The other three panels, from top left to bottom right, display radial profiles of radial velocity, density, and temperature, respectively. Diamonds indicate regions of cold gas with a temperature less than $5 \times 10^4$ K, which we associate with Na i-absorbing gas.

Figure 11 presents a similar analysis but with the highest resolution of 0.1 pc. The radial profiles in all panels use the horizontal axis labeled on the bottom right.

Ejects nearly all its kinetic energy in the hot gas (Mac Low & Ferrara 1999), distributed supernovas deposit as much as half of the energy in the cold gas. However, much of this energy is radiated promptly rather than converted to kinetic energy, so determining the significance of distributed supernovas requires further work.
Figure 12. Column density distributions of cool gas at sightlines through the center as a function of angle extended from the vertical axis in models with 0.2 pc (solid line; model X1) and 0.1 pc (dashed line; X1-0) resolution. The observed range of column density inferred from observations of Na i absorption profiles is $1.0 \times 10^{19}$–$4.3 \times 10^{21} \text{ cm}^{-2}$, shown in dotted lines. Note that the column densities from the models will be reduced over time due to spherical expansion.

5.2.1. Cool Gas Column Densities

Figure 7 shows that the lines of sight plotted in Figures 10 and 11 go through multiple distinct shells, adding to a total column density of cool gas $N_H \sim 7 \times 10^{21} \text{ cm}^{-2}$. The inferred column densities in ultraluminous starbursts are a bit lower, exceeding $10^{21} \text{ cm}^{-2}$ in only one out of four ultraluminous starbursts (Martin 2005). Since the model column densities fluctuate by as much as an order of magnitude between neighboring sightlines, we plot the column density distribution as a function of angle in Figure 12. We find that most sightlines with $N_H > 5 \times 10^{22} \text{ cm}^{-2}$ lie within 30° of the galactic disk where the shells are usually not subject to hydrodynamic instabilities. The typical value is $N_H \sim 10^{22} \text{ cm}^{-2}$, similar to the largest column density estimated from observations of ultraluminous starbursts (Martin 2005, 2006; Rupke et al. 2002).

5.2.2. Geometric Dilution

The column densities found in our model are measured at a time very early in the lifetime of the wind, when it has not expanded far away from the galaxy. However, these winds will typically be observed at later times, when the wind has expanded further. Geometric dilution will then reduce the column density along any particular line of sight even if the gas simply expands radially outward on ballistic orbits. For spherical geometry (of any opening angle) and constant velocity flow, the volume-averaged gas density declines with increasing radius. The amount of dilution is therefore determined by the radial advance of the innermost cold gas. At early times, see Figure 7 this innermost cold gas lies at a radius of $\sim 100 \text{ pc}$, near the wind termination shock. Assuming the cold gas directly above the starburst region flows upward on a ballistic path, we would expect it to move outward by a factor of roughly 100 in the next 25 Myr, reducing the cold gas column density by the same factor.

5.2.3. Resolution Effects

A second reason we overestimate the cool gas column density is that the fragmentation of the shells is limited by numerical resolution, as well as our assumption of azimuthal symmetry. However, as discussed above in Section 4.2.1, the overall kinematics described in Section 5.1 should remain similar unless shell fragments are entirely mixed with hot gas and destroyed. To distinguish real physical effects from numerical artifacts, we study the kinematics of the cold gas along a $\theta = 190°$ line of sight in our fiducial simulation X1 and the high-resolution simulation X1-0.

The first two fragments from the center are parts of a filament remaining from the initial fragmentation of the swept-up supershell due to R–T instability at $t \sim 0.1 \text{ Myr}$. These high-density peaks are traveling with similar velocities $v \sim 400$–$500 \text{ km s}^{-1}$ in both simulations and dominate the column densities of cool gas in the studied sightlines. The outermost shell in X1 keeps sweeping up high-altitude disk gas, while in X1-0 this shell further fragments by R–T instability to the last two fragments in Figure 11. The coherence of the shell in X1 clearly occurs because of the lower resolution in that run. However, the shell in X1 and the outermost fragment of the two in X1-0 travel with similarly high velocity of $\sim 800$–$900 \text{ km s}^{-1}$, though they contribute very little to the total column density.

The hot wind overruns and shocks the first fragment in X1 and the first two fragments in X1-0. As a result, secondary R–T and K–H instabilities act on them, removing gas from the cold cloud and mixing it into the hot wind. In both cases, these clumps of gas are resolved by about $\sim 10$ cells, demonstrating that this is a minimum size below which fragments tend to survive because the secondary instabilities cannot be resolved. Mac Low & Zahnle (1994) show that fragments must be resolved by at least 25 zones to resolve the secondary instabilities. Fragments smaller than that remain as artificially massive clouds in our simulations, overpredicting the cool gas column densities. Klein et al. (1994) suggest that these clumps of gas should be destroyed, via hydrodynamic instabilities, over $\sim 10$ shock crossing timescales

$$t_{cc} = r_c(v_w - v_c)\sqrt{\rho_c/\rho_w} = (0.04 \text{ Myr})\left(\frac{r_c}{2 \text{ pc}}\right) \times \left(\frac{v_w - v_c}{500 \text{ km s}^{-1}}\right)\left(\frac{\rho_c/\rho_w}{100}\right)^{1/2},$$

where $r_c$ is the radius of the clump, $\rho_w$ and $\rho_c$ are the density, and $v_w$ and $v_c$ are the velocity of the wind and the clump, and the scaling parameters hold for typical 10 zone clumps in our model. In reality, $t_{cc}$ may be substantially longer because the density ratios $\rho_c/\rho_w$ are probably underestimated: shell densities, and thus clump densities, are underresolved, while the wind density may be overestimated by our mass-loading scheme. We need $\rho_c/\rho_w \approx 10^4$ for these clumps to survive for more than $\sim 5 \text{ Myr}$, which may just be reachable in a rapidly diverging wind.

Moreover, recent numerical studies clumps may be stabilized against Kelvin–Helmholtz and R–T instabilities, reducing mass loss, by either thermal conduction (Marcolini et al. 2005; Vieser & Hensler 2007a), or even weak magnetic fields (Mac Low et al. 1994; Shin et al. 2008). Since we are not able to address the fate of shell fragments further in this study, we assume the bulk of the cool mass remains cool for the duration of starburst to be observed as Na i-absorbing gas. Larger, denser clumps are more likely to survive, in reality.
5.3. Velocity Spread

The bottom left panels of Figures 10 and 11 show the distribution of column density of cool gas $N_H$ as a function of its radial velocity $v_r$ at the end of the simulations. So long as the linear version of the ballistic approximation is correct, the velocity spread remains constant. Cool gas is seen over a wide range of velocities > 450 km s$^{-1}$ at both resolutions. Shells that fragment earlier have been accelerated less, and so their fragments travel more slowly.

We now explore whether absorption from multiple shell fragments can explain the large absorption line widths measured for ultraluminous starbursts by considering the effects of resolution and viewing angle. Figure 13 shows the distribution of velocity widths $\Delta v$ seen in cool gas with $T < 5 \times 10^4$ K in runs X1-0, X1, X1-2, and X1-4, with grid size increasing from 0.1 to 0.8 pc, along sightlines spaced by 1° from the vertical axis. We define $\Delta v$ as the difference between the maximum and the minimum velocities seen in the cool gas along a given sightline—that is, full width at zero. This resolution study demonstrates that the fraction of sightlines with large velocity spread increases as the resolution improves up to 0.2 pc, but then appears to converge.

Comparison to Figure 7 shows that the largest velocity widths are seen in sightlines intersecting the largest number of fragmented shells. Better resolving postshock densities in the swept-up, cooled, shells leads to quicker shell fragmentation and reforming, producing a larger number of shells, and thus the wider range of velocities in the fragments.

Figure 13 also shows that most sightlines in X1 and X1-0 are above the average observed line width $(\langle v \rangle = 320 \pm 120$ km s$^{-1}$) (dashed lines). Observers measure the average of the sightlines toward continuum sources subtending about an arcsecond, and this corresponds to a length of 1.84 $h_{0.7}$ kpc at redshift 0.1. Assuming that the observers are likely to view ULIRGs at $t = 5–10$ Myr after the onset of starburst, we can suggest all the shell fragments within $\sim20°–40°$ at a given sightline in our simulations will contribute to the absorption profiles. Hence, our models suggest that observers will measure a large line width regardless of viewing angle.

To quantify this, we compute the average velocity width of all sightlines in an axisymmetric cone within $\theta$ as

$$\bar{\Delta v}(\le \theta) = \frac{\int_0^\theta \Delta v(\alpha) \sin(\alpha) d\alpha}{\int_0^\theta \sin(\alpha) d\alpha}.$$  

The value of $\bar{\Delta v}$ is lowered by sightlines near the disk midplane where the bubbles are not blowing out and sightlines intersecting holes made by blowout with nearly zero $\Delta v$. The average velocity widths are clearly smaller for lower resolution runs, however. We find $\Delta v(\theta \le 30°) = 220$ km s$^{-1}$ for run X1-0, and 210 km s$^{-1}$ for X1, but only $110$ km s$^{-1}$ for X1-4.

5.4. Average Velocity

Figure 14 shows the distribution as a function of angle with degree spacing of mass-weighted average velocities of cool gas $v_{\text{av}}$ (solid lines) for our standard resolution study (runs X1-0 through X1-4). The average shell velocities plotted may be misleading, since shell mass differs significantly at each sightline. Therefore, we also plot a mass-weighted average shell velocity across a 10° arc centered on each sightline $v_{\text{av}}$. This ought to be the quantity most directly comparable to observations of the average velocity.

On sightlines with multiple fragments, the average mass-weighted velocity reflects the velocity of the more massive fragments. These are fragments of the initial swept up shell,
which follow ballistic orbits with little acceleration after shell fragmentation. Thus, average mass-weighted velocity tends to lie substantially below peak velocity.

The observable quantity $v_{\text{av}, 10}$ shows a clear converging trend with resolution. The converged value toward the pole appears to be under 400 km s$^{-1}$, with correspondingly lower values at other angles, as shown in Figure 14. Figure 14 suggests that cool gas in shells and their fragments will be observed to be traveling with $v_{\text{av}, 10} \approx 200–350$ km s$^{-1}$ at angles to the pole $\theta < 60^\circ$, the angle where blowout occurs in all models.

Our model is consistent with observations that constrain ULIRG wind geometry. Figure 14 shows that absorption at velocities $v > 100$ km s$^{-1}$ is detected at all angles greater than $10^\circ$ from the disk plane. It follows that 98% of all random, radial sightlines would exhibit such outflows. This result is consistent with 15 of 18 observed objects showing such outflows in Martin (2005), and a similar fraction seen by Rupke et al. (2005).

Figure 15 shows the distributions of $v_{\text{av}}$ and $v_{\text{av}, 10}$ for models X2 and X3 run with decreasing mechanical luminosity at the same resolution as model X1. The mass-weighted average velocities of shells should depend on the mechanical luminosity $L_{\text{mech}}$ driving the bubble by thermal pressure. In a spherical bubble, the dependence would be $L_{\text{mech}}^{-1/5}$. This dependence is actually slightly stronger in our models of blowout. We find $v_{\text{av}, 10}$ dropping from 400 km s$^{-1}$ to $\sim 100$ km s$^{-1}$ moving from model X1 to X3, with each step having an order of magnitude lower mechanical luminosity. This gives an empirical relation closer to $L_{\text{mech}}^{-0.3}$. Note the high spike in $v_{\text{av}, 10}$ of X3 around $\sim 15^\circ$ is biased by a small amount of cool mass present in its vicinity. By way of comparison with observations, Figure 6 of Martin (2005) shows velocities reaching $v = 700$ km s$^{-1}$ for an SFR of $1000 M_{\odot}$ yr$^{-1}$. At an SFR of $1 M_{\odot}$ yr$^{-1}$, the sparse data shown in that Figure suggest $v = 30–40$ km s$^{-1}$, while our empirical relation would suggest a value $v \approx 90$ km s$^{-1}$. More data at low SFRs and a broader range of models will be required to establish whether there is truly a discrepancy with the model results.

The mass-weighted average velocities $v_{\text{av}, 10}$ in X1 and X2 agree with the observed shell velocity, but $v_{\text{av}, 10}$ in X3 is well below the observed value. In disks as massive as these, starbursts with mechanical luminosity lower than $10^{42}$ erg s$^{-1}$ produce the lower end of the outflow velocity range observed in ULIRGs, which average $v_{\text{obs}} = 330 \pm 100$ km s$^{-1}$ in the Martin (2005) sample and 170 km s$^{-1}$ in the column-density weighted average velocities from the Rupke et al. (2005) sample. Column-density weighting with velocity can be confounded by variation in covering factor with velocity, as has been shown for AGN outflows (e.g., Arav et al. 1999, 2005). We cannot model covering factors well with our two-dimensional simulations, so we postpone consideration of the question of column-weighted velocity to future work. We do note that overpredicting acceleration is much harder to do than underpredicting it, though, so we think our model is likely to be robust. The velocity spread, which is our most important result from the simulations, is similar to that found in both ULIRG studies.

5.5. Terminal Velocity

5.5.1. Resolution

Figure 14 also shows the distribution of terminal velocity of cool gas $v_{\text{term}}$ at each sightline. We define the maximum velocity at blowout as the terminal velocity a shell will ever acquire, following the ballistic approximation (Section 5.1). The maximum velocity of shell fragments is very high; $v_{\text{term}} > 500$ km s$^{-1}$ at the angle $\theta < 60^\circ$ where blowout occurs in all the runs. This high-velocity cool gas is found in the fragments of the outermost shells in our highest resolution model X1-0. The hot wind continues to accelerate a piece of shell, sweeping up the ambient gas, until it fragments further. As discussed above, shell fragmentation is very sensitive to resolution. Thus it is important to test the convergence of the mass of high-velocity gas.

Figure 16(a) shows the mass distribution of cool gas as a function of velocity in X1-0, X1, and X1-2 at $t = 0.27$ Myr. The fraction of cool gas traveling at high velocities is low. For example, the mass traveling with $v \geq 500$ km s$^{-1}$ is less than a few percent of the total shell mass, and the mass traveling faster than the observed terminal velocity of 750 km s$^{-1}$ is $< 0.1\%$. However, these results are best understood as upper limits. The amounts of cool gas in the high-velocity end are progressively smaller as the resolution increases. Roughly factor of 2 decreases in the mass of cool gas with velocities above 500 km s$^{-1}$ occur between runs with factor of 2 improvements in linear resolution. This suggests that we have not yet converged on the actual mass of high-velocity gas, although we have set good upper limits. We think this high-velocity gas is not likely to go away entirely, even if we run a simulation with a higher resolution and with additional physics. However, we cannot fully quantify its amount with our simulations.

The convergence properties of the peak velocity for cold gas $v_{\text{term}}$ are somewhat better. The highest resolution models show $< 20\%$ variations, suggesting that the general result is reasonably robust. The observed maximum value of 750 km s$^{-1}$ is consistent with the models at all $\theta < 60^\circ$.

5.5.2. Mass Loading

Figure 17 shows the distributions of $v_{\text{av}}$, $v_{\text{av}, 10}$, and $v_{\text{term}}$ for simulations with different mass-loading rates and so different wind terminal velocities: U1-A, U1, U1-B, and U1-C. Cool gas has higher terminal velocities in bubbles with hotter winds that themselves have higher velocities $v_{\text{wind}}$. The hottest wind, in model U1-A, has $v_{\text{wind}} = 2700$ km s$^{-1}$, while the coldest wind, in U1-C, has $v_{\text{wind}} = 250$ km s$^{-1}$. Figure 16(b) shows the mass distribution of cool gas as a function of velocity for the four runs at the time of blowout. The cool gas with high $v_{\text{term}}$ carries only a small amount of mass (see Section 5.5.1). For example, the fraction of cool gas with $v > 500$ km s$^{-1}$ is $\lesssim 2\%$ of the total shell mass, and the fraction with $v > 750$ km s$^{-1}$ is $\lesssim 0.5\%$. The bulk of swept-up and cooled gas is driven to $v \sim 400$ km s$^{-1}$.
by thermal pressure of hot sonic gas, but a small fraction of it seems to be accelerated to higher velocity by the ram pressure of the same hot gas as it accelerates to supersonic velocities during blowout.

5.5.3. Mechanical Luminosity

The terminal wind velocity (Equation (3)) in our models X2 and X3 with lower $L_{\text{mech}}$ is the same as that of our fiducial model X1, because the ejected mass $M_{SN} \propto L_{\text{mech}}$, and we keep the mass-loading factor $\xi$ constant. In all three models $v_{\text{wind}} \approx 1000 \text{ km s}^{-1}$.

Figure 15 shows that the terminal velocities of cool gas are not as high as the observed average terminal velocity, $v_{t,\text{obs}} = 750 \text{ km s}^{-1}$ at most sightlines if $L_{\text{mech}} \leq 10^{42} \text{ erg s}^{-1}$. It is much harder for the wind ram pressure to accelerate the outermost shells and their fragments to very high velocity, $> 500 \text{ km s}^{-1}$ if their starting velocity at blowout $v_b \lesssim 200 \text{ km s}^{-1}$.
5.5.4. Summary

In summary, the terminal velocity of cool gas is determined by the combination of total mechanical power $L_{\text{mech}}$ and wind terminal velocity $v_{\text{wind}}$ (determined by the mass-loading rate $\xi$ in our model). Bubble thermal pressure accelerates the bulk of shells to their final velocity before blowout and ram pressure of the hot transonic wind after blowout accelerates a small fraction of the cool gas to much higher velocities. Faster, less mass-loaded winds accelerate cool gas to higher velocities. Our wind is an energy-driven wind, not a momentum-driven wind (Murray et al. 2005). The thermal energy of the hot wind is gradually turning into its kinetic energy by blowout (recall radiative cooling is turned off inside the wind; see Section 3). Our model accounts for the observed 750 km s$^{-1}$ terminal velocity of Na i-absorbing gas without invoking any additional physics such as radiation pressure.

5.6. Absorption Profiles

To directly compare to observed profiles, we generate Na i $\lambda$5890 absorption line profiles along sightlines through our simulations, as well as generalizing this procedure to fully model the observed doublet Na i $\lambda$, $\lambda$5890, 5896. To generate the profiles, we begin with the line intensity

$$I_v = I_v(0) \exp^{-\tau_i} = I_v(0) \prod_{i=1}^{N} \exp^{-\tau_i} \quad (8)$$

where the optical depth through cell $i$ at frequency $v$ is $\tau_i$, the sightline intersects $N$ cells, and the background continuum is $I_v(0)$. We normalize the continuum by setting $I_v(0) = 1$ and compute the profile as a function of the macroscopic velocity of Na i-absorbing gas, $v$. We compute $I_v$ in the simulations by setting $v = c(v - v_0)/v_0$, and computing the optical depth contributed by each cell in each of 1000 velocity bins. The optical depth in each cell

$$\tau_i(v) = N_{\text{Na}i} \lambda \lambda \text{Na}i P(\Delta v), \quad (9)$$

(Spitzer 1978) where $N_{\text{Na}i}$ is the column density of Na i in a cell $i$, and the absorption cross section $\sigma_i$, integrated over frequency $v$ for $hv \gg kT$ is $s \approx 2.654 \times 10^{-2} f_{5890}$ with the oscillator strength $f_{5890} = 0.6$. The Maxwellian velocity distribution function, $P(\Delta v)$, is

$$P(\Delta v) = \frac{1}{\sqrt{\pi b}} \exp - (\Delta v^2/b^2), \quad (10)$$

with $\Delta v = v - v_i$ and $b = \sqrt{2k_B T_i/\mu m_H}$ for thermal broadening. Our simulations do not track chemical abundances, ionization state, or dust depletion, so we do not directly predict the Na i column. For the purpose of illustration, we compute the total Na i column from the total H i gas column using the conversion used by Martin (2005) to estimate $N_{\text{HI}}$ from their $N_{\text{Na}i}$ measurements, $N_{\text{Na}i} = 1.122 \times 10^{-5} N_{\text{HI}}$.

We generated Na i 5890 line profiles along lines of sight through models X1 and X1-0 to show the effect of numerical resolution on the line profiles in Figure 18. Each sightline is described by its inclination from the polar axis of the simulation. The profile shapes reflect the different structure in these simulations. We can, for example, compare the absorption profiles in the middle left panel in Figure 18 with the density and velocity distributions of shells in Figures 10 and 11. The sightlines at $\theta = 19^\circ$ intersect three and four shell fragments in X1 and X1-0, respectively. Each of these fragments generates a $\Delta v \gtrsim 50–100$ km s$^{-1}$ absorption line. A few of these lines are optically thick with the line profile becoming completely black at the line center. As the wind expands, column densities will drop linearly with radius due to geometric dilution. The complex of lines is spread out over $\sim$400 km s$^{-1}$.

Figure 18 shows that the line profiles from the high and low resolution models present very similar structure for the most part. One minor exception is the minimum velocity. From $10^\circ$ to $20^\circ$, the line profiles show a sharp cutoff at a velocity $\sim 500$ km s$^{-1}$ in X1, but $\sim 3–40$ km s$^{-1}$ in X1-0. This cutoff reflects the velocity of the first R–T fragments that form from the accelerating swept-up shells, which is lower in the high-resolution simulation. The overall lineshape does appear reasonably well converged.

Observers see the average over parallel sightlines subtending a few kiloparsecs of the disk. They must also contend with the instrumental response function, and the doublet nature of the Na i line, which has components at 5890 Å and 5896 Å, with optical depths differing by a factor of $\sim 2$. To compare to the actual observed profiles, we generate the Na i 5890/5896 doublet in the frame of Na i 5890 by assuming the ratio of equivalent widths is only 1.3, typical of the observed lines in ULIRGs (rather than the optically thin limit of two).

We present average line profiles of the doublet over a $20^\circ$ wide set of radial sightlines spaced at degree intervals. Although this is not exactly what is measured, it is comparable because these rays will subtend about 1–2 kpc of the shell when the bubble is $\sim 10$ Myr old, so they intersect the same region of the shell, albeit at slightly different angles. We are actually doing the analysis at an earlier time, when the bubble is a factor of 100 smaller, and correspondingly higher column density, though, so we also must apply a geometric dilution factor to the column densities.
Figure 19. Simulated Na I 5890/5896 doublet absorption profiles averaged over 20° centered on the given (nonuniformly distributed) angles for model X1-0 with \( dx = 0.1 \) pc (solid line). The same profiles are shown after geometric dilution of the column densities by a factor of 100 (dashed line), and application of a Gaussian instrumental broadening with FWHM of 65 km s\(^{-1}\), for comparison with five observed ULIRG spectra (thin solid lines) from Martin (2005). Note that the velocity frame is centered on the 5890 line; blue-shifted absorption from the 5896 line lies at low velocities in this frame. We also show the observed range of average shell velocity in shaded area and the observed average terminal velocity in dash-dot-dot line.

Figure 20. Velocity width is plotted as in Figure 13 (top panels), and the mass-weighted average velocity and the terminal velocity are plotted as in Figure 14 (bottom panels) for model X1. Parallel sightlines are chosen along a slit oriented at \( \theta = 30° \) from the axisymmetric axis and the major axis (left panels) and the minor axis (right panels).

Figure 19 compares our simulated doublet line profiles, with and without geometric dilution of the column densities by a factor of 100, to five typical ULIRG spectra from Martin (2005) that have an instrumental broadening with FWHM \( \simeq \) 65 km s\(^{-1}\). We convolve in a Gaussian of that width to simulate the broadening. The widths of the absorption profiles are very similar between our model at \( \theta \leq 50° \) where blowout occurs and the observations. The undiluted absorption in our model much exceeds that observed. Our model overpredicts the cool gas column for two reasons (see Section 5.2). First, because we analyzed the models at a very early time in the starburst. At later times, the spherical expansion of the wind and the cold gas embedded within it dilutes their column densities. Second, the fragmentation of the shell is incompletely resolved, so some material remains cool that in reality would have been mixed into the hot wind. We account for the first factor by directly reducing the column densities, producing the diluted profiles shown in the figure, which reproduce the observed intensities well.

The ULIRG spectra show absorption beginning from the systemic velocity, and extending to high velocities. Although this zero velocity absorption is absent in Figure 18, it is seen in the simulated doublet profiles shown in Figure 19, as it comes primarily from contributions from the \( \lambda 5896 \) line of the doublet blueshifted into zero velocity of the \( \lambda 5890 \) line (note our model does not include the contribution from stellar absorption).

5.6.1. Long Slit Observations

Measured Doppler shifts across ULIRGs show the cool outflow is extended spatially; and some outflows present a significant velocity gradient over kpc scales (Martin 2006). Our simulation ends long before the wind has reached such large scales and does not include the rotation of the galactic disk. The simulation does demonstrate, however, the amplitude of the velocity gradient that arises across the minor axis due to projection effects.

The wind is launched without any net angular momentum in our two-dimensional simulation. We examined whether, in the absence of rotation, the position affected the velocity much. Figure 20 shows terminal and mass-weighted average velocity along slits oriented at 30° from the major and minor axes of the disk, using parallel lines of sight through a three-dimensional simulation. The absorption properties along the major axis of the galaxy will be symmetric about its center. An asymmetric velocity gradient is produced along the minor axis because one side of the outflow cone is directed more along the sightline. Substantial variations in velocity width and average velocity are seen at \( \pm 40 \) pc and \( \pm 30 \) pc. These occur where sightlines intersect massive, slow-moving shells at \( \theta > 45° \) (see Figure 14) as well as light, fast-moving blowout fragments at \( \theta < 45° \). On the other hand, very high terminal velocities are seen on most sightlines since they intersect the fast blowout components.

The simulated sightlines encounter both fast-moving and slow-moving gas at various latitudes, unlike the radial sightlines that we studied in previous sections. Our simulations viewed along nonradial sightlines naturally account for a wide velocity range of cool gas starting at \( v \approx 0 \) and a high terminal velocity.

5.7. Line Widths of Na I-Absorbing and Hα-Emitting Gas

Finally, we crudely try to separate low-ionization, Na I-absorbing gas and ionized Hα-emitting gas in the cooled shell fragments in our simulations. Our aim is to compare the line widths of both components. Heckman et al. (2000) observed no correlation between the absorption and emission line widths for these lines, and drew the conclusion that only one of the lines could originate in the swept-up shells.
To model the Hα emission, we must decide whether the emitting gas is primarily photoionized by the central starburst or shock ionized in the wind. Shock ionization dominates at least some starburst-driven winds, such as NGC 1482 (Veilleux & Rupke 2002). Some observed ULIRGs also show extended shock-excited nebulae (Monreal-Ibero et al. 2006). However, the dynamics observed by Heckman et al. (2000) seem unlikely to come from shock ionized gas, since they see no correlation with the motions of the cold gas that presumably trace the shock. Therefore we assume photoionization by the central starburst, choosing an ionizing photon luminosity $Q$ and a density distribution from a time the galaxy is likely to be observed. Figure 2 shows the ionizing photon luminosities as a function of time for $Q = 10^{53}, 10^{54}, 1.9 \times 10^{54}$, and $10^{55}$ photons s$^{-1}$ in a line of sight through the center at an angle of $13^\circ$ from the vertical axis in model U1. The velocity widths of the two components vary as $Q$ is changed.

To model the Hα emission, we must decide whether the emitting gas is primarily photoionized by the central starburst or shock ionized in the wind. Shock ionization dominates at least some starburst-driven winds, such as NGC 1482 (Veilleux & Rupke 2002). Some observed ULIRGs also show extended shock-excited nebulae (Monreal-Ibero et al. 2006). However, the dynamics observed by Heckman et al. (2000) seem unlikely to come from shock ionized gas, since they see no correlation with the motions of the cold gas that presumably trace the shock. Therefore we assume photoionization by the central starburst, choosing an ionizing photon luminosity $Q$ and a density distribution from a time the galaxy is likely to be observed. Figure 2 shows the ionizing photon luminosities as a function of time for three starburst cases that we consider. Since the propagation of ionization fronts critically depends on both the densities and the positions of shells, we cannot extrapolate the density distribution at blowout to $t \approx 10$ Myr using the ballistic approximation. We can, however, still vary the strength of photon luminosity over a few orders of magnitude to examine the effect of attenuating the photon flux by $1/r^2$ on our existing simulations. For example, a shell will experience about three orders of magnitude less photons per unit area after it travels with 500 km s$^{-1}$ for 10 Myr from $r_p = 200$ pc. For our purpose of demonstrating the lack of correlation between the line width of Na i-absorbing and Hα-emitting gas, this crude method is sufficient.

To solve for the transfer of ionizing radiation across our grids, we use the photon-conserving radiative transfer code developed by Abel et al. (1999) in the same manner as it was used in Fujita et al. (2003). It computes the propagation of ionization fronts around a point source; in our study this is a central starburst source. This code is a post-processing step that operates on a given density distribution at a given time step in our simulation. However, the ionization fronts propagate sufficiently rapidly to adjust almost instantaneously to a changing density distribution.

Figure 21 shows the column density as a function of velocity for Na i-absorbing gas (diamonds) and Hα-emitting gas (triangles) in run U1 at $\theta = 13^\circ$ with ionizing luminosities of $Q = 10^{53}, 10^{54}, 1.9 \times 10^{54}$, and $10^{55}$ photons s$^{-1}$. With $Q \lesssim 10^{54}$ photons s$^{-1}$, the line widths are very small for Hα-emitting gas, $\Delta v_{\text{Hα}} < 100$ km s$^{-1}$, but large for Na i-absorbing gas, $\Delta v_{\text{Na}} \approx 480$ km s$^{-1}$. The densest shell at $r = 110$ pc traps the photons rather effectively. As $Q$ is increased, that first shell is ionized, but the second shell then traps the photons. The bottom left panel shows that the Na i line width is now very small, $\Delta v_{\text{Na}} = 38$ km s$^{-1}$, while $\Delta v_{\text{Hα}} = 330$ km s$^{-1}$. With $Q \gtrsim 2 \times 10^{54}$, all the shells are ionized. In reality, the transition from large $\Delta v_{\text{Na}}$ to $\Delta v_{\text{Na}} \approx 0$ should not be this abrupt because many combinations of fragments and clumps are possible within the observers’ $10 \times 10^2$ field of view. As we note, we are far from presenting realistic distributions of Na i-absorbing and Hα-emitting gas. However, Figure 21 still demonstrates that the line widths of Na i-absorbing and Hα-emitting gas can easily be uncorrelated although they both originate in the swept-up shells.

6. CAVEATS AND SUMMARY

6.1. Caveats

This study clearly is not the final word on this subject. Rather it is a proof of the concept that a starburst wind can accelerate neutral gas up to high velocities without special circumstances, and that the resulting flow configuration can reproduce many of the observable properties of ULIRG winds. We here recap the approximations we have made in order to treat this problem, and add some discussion of how they might be lifted.

The biggest approximation of our study is that we analyze the kinematics of shell fragments at $t \ll 1$ Myr in our small grids and extrapolate the results by the ballistic approximation including geometric dilution for comparison with the observations at $t \gg 1$ Myr. Much of the total mechanical energy from even an instantaneous burst of star formation has yet to be deposited at an age of 2 Myr when our simulation ends. These dense clumps may further fragment and ultimately be destroyed by R–T and Kelvin–Helmholtz instability as the high-velocity flow of gas from within the bubble streams through them. Conversely, they may be dense enough to be accelerated further prior to their destruction by the ongoing starburst wind.

As the grid resolution of our models improves to 0.1 pc, we do begin to resolve the details of hydrodynamic instabilities. Marcolini et al. (2005) used 0.1 pc resolution for their simulations of the dynamical shredding of radiatively cooling clouds, and seem to have reached adequate resolution. Thus, in future work, continuing to evolve our bubbles in bigger grids with the same resolution may well improve our results.

We have made four other substantial physical approximations as well. First, the assumption of azimuthal symmetry further suppresses R–T instability. Mac Low et al. (1989) pointed out that the typical spike and bubble structure of the R–T instability is limited to rings in an axisymmetric blowout. The detailed distribution of the fragments will certainly be different in three dimensions, but finer fragments will probably have a broader velocity range, and thus are unlikely to change our qualitative result. The behavior of individual fragments will also not be qualitatively changed by the dimensionality, although two-dimensional fragments split into larger pieces initially (Stone & Norman 1992b; Korycansky et al. 2002).

Second, we neglected magnetic fields in this study. We showed that magnetic fields reduce the peak shell density, and thus the fragmentation (Section 4.2.1), but that they don’t act to do so before fragmentation is essentially complete in our
model, so that this effect is secondary. Magnetic fields can, on the other hand, help preserve individual fragments from further breakup (Stone & Norman 1992b; Mac Low et al. 1994; Shin et al. 2008), possibly even allowing their further acceleration in the continuing starburst wind.

Third, thermal conduction is not explicitly treated in our model. This acts on parsec scales, so numerical diffusion and turbulent mixing will still dominate over thermal conduction under most circumstances at the resolutions that we consider. It is worth noting, though, that studies of individual clump fragmentation have found that thermal conduction can have stabilizing effects (Marcolini et al. 2005; Vieser & Hensler 2007b). Fourth, the dynamical effects of photoionization have also been neglected. This could heat at least low column-density shells and fragments up to $10^4$ K, reducing their density contrast with the wind and enhancing their tendency to fragment.

We also assume a single starburst at the disk center for modeling a galactic outflow. In realistic galaxies, we expect multiple star clusters to be scattered around the disks (e.g., Vacca 1996; Martin 1998). A more complicated structure of cool shells and their fragments is expected with a realistic distribution of star formation as shown by models of supernova-driven turbulence in our own galaxy (e.g., Avillez 2000; Avillez & Berry 2001; Joung & Mac Low 2006). Such a distribution was modeled with a fractal density distribution by Cooper et al. (2008). Fragile et al. (2004) showed that in dwarf galaxies, distributed supernova explosions resulted in increased transfer of energy to the cold gas compared to the centrally concentrated energy injection assumed here and by Mac Low & Ferrara (1999), so our approximation likely gives a lower limit to the kinetic energy of the cold gas. This supports our result of a wide velocity range in Na I-absorbing gas. The next obvious step is to study the effects of realistic star formation on the shell kinematics in three-dimensional, AMR simulations.

Another related issue is that our models assume a low-density, uniform gas distribution above the site of blowout. However, ULIRGs are formed in galaxy mergers. The actual situation near the nucleus of merging galaxies is likely to be more chaotic. We rely on the idea that the large-scale tidal tails and other merger structures are generally going to be well removed from the wind generation region and will not cover a large fraction of the sky. Ultimately, to remove this limitation, full models of merging galaxies at high resolution will be needed, but this remains some years in the future.

6.2. Summary

We study the origin and the kinematics of cool gas that produces Na I absorption lines in galactic winds, using hydrodynamic simulations of the blowout of starburst-driven superbubbles from the molecular disk of ULIRGs. The bubbles sweep up the dense disk gas, which quickly cools to form dense, thin shells. The cooled shells fragment by R–T instability as they accelerate in the stratified atmosphere of the disk. This blowout occurs very early in our models, at $t \ll 1$ Myr in models with $L_{\text{mech}} \gtrsim 10^{41}$ erg s$^{-1}$. The dense fragments left by the R–T instability lag behind the low-density, high-velocity wind. These fragments carry most of the mass that is swept up by the bubbles, and should have the highest column of Na I. The results of our numerical convergence study suggest that superbubble blowout, combined with subsequent geometrical dilution, can reproduce not just qualitative but quantitative properties of the observed lines.
As a result of R–T fragmentation, multiple shell fragments and clumps travel at different velocities. A sightline going through them reproduces the observed broad line width of the Na i absorption profiles: \( \langle v \rangle = 320 \pm 120 \text{ km s}^{-1} \). This result does not appear to depend strongly on physical parameters such as mass-loading rate, mechanical luminosity, or surface density. However, this result requires sufficient numerical resolution to follow secondary fragmentation of the shell, so that any line of sight runs over multiple cold gas fragments. We find that a resolution of at least \( \sim 0.2 \text{ pc} \) is required to produce this effect, if mass-loading rates remain moderate \( \xi \lesssim 10 \). The suggestion by Heckman et al. (2000) that swept-up shells will not show a high velocity is less than a few percent of that of the total interior gas, which is proportional to \( L > L_{\text{mech}} \). No other parameters influence the results. The mass-weighted average velocities in our high-resolution simulation with \( L_{\text{mech}} = 10^{43} \text{ erg s}^{-1} \) agree with the observed value \( v_{\text{obs}} = 330 \pm 100 \text{ km s}^{-1} \). A mechanical luminosity of \( L_{\text{mech}} = 10^{43} \text{ erg s}^{-1} \) corresponds to an instantaneous burst of \( M_* = 10^3 \text{ M}_\odot \) or a continuous SFR= 500 \text{ M}_\odot \text{ yr}^{-1} \) at \( t < 1 \text{ Myr} \) after the onset of starburst.

As the swept-up shells fragment by R–T instability, the bubble interior gas blows out and becomes a high-density supersonic wind, as its thermal energy is converted to kinetic energy by expansion. The ram pressure of this wind continues to accelerate entrained cold fragments, and to sweep up high-velocity gas. These outermost shells and fragments with high velocity do fragment further in our highest-resolution simulation, but their fragments are traveling with similarly high velocity. Although the mass in the high-velocity tail decreased as the resolution is increased, we believe it is not likely to go away entirely in a higher-resolution simulation.

The clumps and fragments seen in the simulations may be observed as Na i-absorbing gas or \( \text{H}\alpha \)-emitting gas, depending on the amount of ionizing photons produced from the central starburst source. To study this, we used a ray-tracing method to model the location of the ionization front in our models as a function of time, and so to trace the gas emitting in \( \text{H}\alpha \). By varying the photon luminosity, we showed that the velocity range of the ionized and neutral components do not show any correlation with each other. Thus the lack of correlation in the observed line widths of Na i absorbing gas and \( \text{H}\alpha \)-emitting gas does not rule out the swept-up shells as the origin of both components.

The hot wind in our model is purely energy-driven. By construction the interior bubble can never become momentum-driven because radiative cooling is turned off in the bubble interior in our models. Future work must determine whether this mechanism reproduces the empirical correlation observed between the maximum outflow velocity and the escape velocity of the host galaxy (Martin 2005). It should also predict the scaling of the mass-loss efficiency with galaxy mass, a key input in cosmological models (e.g., Oppenheimer & Davé 2006, 2008).

We thank M. L. Norman and the Laboratory for Computational Astrophysics for use of ZEUS-3D, R. F. Coker, G. Gisler, R. M. Hueckstaedt, and C. Scovel for the help with getting started with SAGE, and M. K. R. Joung for useful discussions. We thank the referee, D. Rupke, for his extensive assistance in placing this work in context. Computations were performed on the SGI Origin 2000 of the Rose Center for Earth and Space, and the QSC machine of Los Alamos National Laboratory (LANL). A.F. was supported by a cooperative agreement between UCSB and LANL. C.L.M. thanks the Alfred P. Sloan Foundation and the David and Lucile Packard Foundation for support for this work. M.-M.M.L. was partly supported by NSF grants AST99-85392 and AST03-07854, and by stipends from the Max-Planck-Gesellschaft and the DAAD.
APPENDIX

BLOWOUT IN A DWARF GALAXY WITH ZEUS AND SAGE

A major issue in our simulations is whether we sufficiently resolve the unstable shell during blowout. To further examine this question, in this appendix we describe a comparison between models of blowout from a dwarf galaxy similar to that described by Fujita et al. (2003) run with ZEUS-3D, and single-grid and AMR versions of SAGE.

SAGE is an AMR hydrodynamic code developed at LANL/SAIC. It is second order accurate using a piecewise linear Godunov scheme (Kerbyson et al. 2001; Gittings et al. 2008).

The dwarf galaxy has a nominal redshift $z = 8$, and disk mass $M_d = 10^8 M_\odot$. We choose the midplane density of this galactic disk so that it has an exponential surface density profile. We set up the gas in hydrostatic equilibrium with a background galactic disk so that it has an exponential surface density profile.

In SAGE, the cooling is solved with an explicit method that subcyles the cooling source terms, based on a radiative cooling routine that solves for nonequilibrium chemistry (Abel et al. 1997). The time steps for each subcycle are determined as $e\epsilon/e$, where $e$ is internal energy density and $\epsilon = 0.1$. In ZEUS-3D, on the other hand, we assume that the cooling rate is a function of temperature only, using calculations of equilibrium ionization cooling rates by Sutherland & Dopita (1993) with a cutoff at $T < 10^4$ K, and ignore inverse Compton cooling for simplicity. The cooling curve is implemented in the energy equation with cooling rates by Sutherland & Dopita (1993) with a cutoff at temperature only, using calculations of equilibrium ionization.

The cooling curve is implemented in the energy equation with cooling rates by Sutherland & Dopita (1993) with a cutoff at temperature only, using calculations of equilibrium ionization. This curve is chosen so that it is diffusive than the Van Leer (1977) method of ZEUS. However, we show in Figure 22 that the SAGE/AMR run begins resolving fine structures caused by R–T instability as well as ZEUS-3D run when one more level of refinement is added, doubling the maximum effective resolution. Even in this case, the SAGE/AMR run still uses a smaller number of total and active cells than the ZEUS-3D run. We can save noticeable computational time with SAGE/AMR as we tackle a problem with more than a few million cells. This advantage will be bigger with three-dimensional models.

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