First Results from SMAUG: Characterization of Multiphase Galactic Outflows from a Suite of Local Star-forming Galactic Disk Simulations

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Abstract

Large-scale outflows in star-forming galaxies are observed to be ubiquitous and are a key aspect of theoretical modeling of galactic evolution, the focus of the Simulating Multiscale Astrophysics to Understand Galaxies (SMAUG) project. Gas blown out from galactic disks, similar to gas within galaxies, consists of multiple phases with large contrasts of density, temperature, and other properties. To study multiphase outflows as emergent phenomena, we run a suite of roughly parsec-resolution local galactic disk simulations using the TIGRESS framework. Explicit modeling of the interstellar medium (ISM), including star formation and self-consistent radiative heating plus supernova feedback, regulates ISM properties and drives the outflow. We investigate the scaling of outflow mass, momentum, energy, and metal loading factors with galactic disk properties, including star formation rate (SFR) surface density (ΣSFR ∼ 10−4 − 1M⊙ kpc−2 yr−1), gas surface density (Σgas ∼ 1–100 M⊙ pc−2), and total midplane pressure (or weight; Pmid ≈ W ∼ 107–109 kR cm−3 K). The main components of outflowing gas are mass-delivering cool gas (T ∼ 104 K) and energy/metal-delivering hot gas (T ∼ 107 K). Cool mass outflow rates measured at outflow launch points (one or two scale heights ∼300 pc–1 kpc) are 1–100 times the SFR (decreasing with ΣSFRG), although in massive galaxies most mass falls back owing to insufficient outflow velocity. The hot galactic outflow carries mass comparable to 10% of the SFR, together with 10%–20% of the energy and 30%–60% of the metal mass injected by SN feedback. Importantly, our analysis demonstrates that in any physically motivated cosmological wind model it is crucial to include at least two distinct thermal wind components.

Unified Astronomy Thesaurus concepts: Galactic winds (572); Magnetohydrodynamical simulations (1966); Star formation (1569); Stellar feedback (1602); Interstellar medium (847)

Supporting material: animations

1. Introduction

In current theories of galaxy formation and evolution, galactic winds are an important element, counteracting cosmic accretion to limit stellar mass growth of galaxies. Even in the earliest theoretical models of galaxy formation in dark matter halos, the issue of overproduction of stellar mass was recognized, necessitating mass and energy flows out of galaxies (e.g., White & Rees 1978; Dekel & Silk 1986; White & Frenk 1991). Recent cosmological hydrodynamical simulations and semianalytic models (SAMs) that successfully match the observed galaxy statistics, including stellar mass–halo mass relations, all require ejection of a significant fraction of the gas mass accreted in the form of galactic-scale winds (see reviews of Somerville & Davé 2015; Naab & Ostriker 2017, and references therein). With the enormous spatial and temporal domains required for cosmological-scale modeling, however, it is not possible to simultaneously represent the detailed properties of the star-forming interstellar medium (ISM) that lead to the production of galactic winds. Absent a means to directly model the physics within galactic disks, the usual practice is to adopt parameterized scaling relations (for both star formation rates (SFRs) and wind mass-loss rates), calibrating free parameters by reference to observations (e.g., Vogelsberger et al. 2013; Crain et al. 2015; Pillepich et al. 2018). The approach of empirically constrained parameterization, while heretofore unavoidable, has been a major source of uncertainty in modern galaxy formation theory. The SMAUG project was initiated to address the need for developing and implementing subgrid treatments for cosmological models that are derived and calibrated from simulations that explicitly model and resolve key physical processes.

In addition to the important role of winds in the theory of galaxy formation, galactic outflows are prevalent in observations of nearby dwarf starbursts and luminous/ultraluminous infrared galaxies (LIRGs/ULIRGs; e.g., Heckman et al. 1990, 2000, 2015; Martin 1999, 2005; Chisholm et al. 2015). Winds appear to be even more ubiquitous in both active galactic nucleus (AGN) host and star-forming galaxies at high redshift (e.g., Pettini et al. 2001; Shapley et al. 2003; Tremonti et al. 2007; Steidel et al. 2010; Erb et al. 2012; Förster Schreiber et al. 2019), although the limited spatial resolution of these observations makes interpretation more difficult.
Galactic outflows driven by star formation are the result of the feedback that is produced by populations of young stars. This feedback—primarily associated with core-collapse supernovae (SNe) from massive stars, but with some contribution from stellar winds and radiation—returns metal-enriched gas at extremely high velocity to the surrounding ISM. As a result of complex interactions driven by SN shocks (and potentially involving cosmic rays as an intermediary), a portion of the ISM gas is accelerated sufficiently to emerge as a galactic wind, delivering mass, momentum, energy, and metals to the circumgalactic/intergalactic medium (CGM/IGM). Because the massive stars responsible for feedback are buried deep within the ISM, properties of galactic outflows are not simply set by the immediate deposition at feedback sites. Instead, localized energy injection events build up expanding bubbles (or superbubbles for correlated feedback events) with more and more momentum as they sweep up surrounding gas (see, e.g., Taylor 1950; Sedov 1959; Cox 1972; McKee & Ostriker 1977, for a single SN, and, e.g., Weaver et al. 1977; McCray & Kafatos 1987; Koo & McKee 1992, for stellar winds or clustered SNe). The hot ISM is produced by stellar-wind and SN shocks and fills the interior of each bubble. Cooling of the shocked ISM when bubble expansion slows to \( \lesssim 200 \text{ km s}^{-1} \) limits the momentum injection for the case of a single SN (e.g., Cioffi et al. 1988; Thornton et al. 1998), while mixing of hot diffuse gas with dense gas at the bubble-shell boundary drains energy from superbubbles and reduces their dynamical impact (e.g., El-Badry et al. 2019). When extreme star formation/feedback events occur, superbubble breakout from the ISM before the onset of cooling alters the dynamics significantly (e.g., Tomisaka & Ikeuchi 1986; Mac Low et al. 1989; Cooper et al. 2008) and enables delivery of a large fraction of pristine metals and original feedback energy to the CGM (e.g., Kim et al. 2017a; Fielding et al. 2018).

Based on focused high-resolution numerical simulations with an inhomogeneous ISM (e.g., Iffrig & Hennebelle 2015; Kim & Ostriker 2015a; Martizzi et al. 2015; Walch & Naab 2015), the net terminal momentum injection from single SNe has been shown to be quite insensitive to the background medium’s average density and detailed structure (this insensitivity is because the onset of cooling by metal lines at postshock temperature \( T \sim 10^{6} \text{ K} \) is insensitive to the density). A practical application of this result to galactic simulations is the numerical approach of injecting the previously calibrated terminal momentum if the energy-conserving stage of SN remnant (SNR) expansion is unresolved (e.g., Hopkins et al. 2014, 2018a; Kimm & Cen 2014; Kim & Ostriker 2017). This “momentum feedback” approach captures the dynamical impact of SN feedback on the warm–cold ISM phases (essentially all of the ISM’s mass) reasonably well, especially for driving turbulence and therefore self-regulating the SFR, even at relatively low numerical resolution (e.g., Kim et al. 2011, 2013; Shetty & Ostriker 2012; Hopkins et al. 2014; Kim & Ostriker 2015b, 2017; Kimm et al. 2015). However, how much hot gas is created in the ISM by expanding SN-driven bubbles, and how much is retained to vent from the ISM into the CGM, depends sensitively on the details of micro/macrophysics and conditions of the vertically stratified ISM.

Physical elements that affect momentum injection and hot gas yield include turbulence, inhomogeneity, magnetization, thermal conduction, and temporal and spatial correlations of feedback (e.g., Kim & Ostriker 2015a; Gentry et al. 2017, 2019; Kim et al. 2017a; Fielding et al. 2018; El-Badry et al. 2019). Although in principle hot gas generation can be implemented via deposition of “residual” thermal energy (e.g., Martizzi et al. 2015), this will be immediately lost if resolution is too low and hot diffuse gas is not spatially separated from warm fast gas (Hu 2019). In general, outflow properties are much more sensitive to resolution than SFRs (e.g., Kim & Ostriker 2017, 2018; Rosdahl et al. 2017; Smith et al. 2018); because proper hot gas generation and evolution is crucial, outflow properties will be incorrect if most SNe are realized in the form of momentum feedback (e.g., Kim & Ostriker 2018; Hu 2019).

A key characteristic of galactic outflows, which is often overlooked in theoretical modeling, is their multiphase nature. Galactic outflows in observations are often detected in spectra of neutral and ionized gas tracers (e.g., Na I, He I, Si II, Si IV; Heckman et al. 1990, 2015; Martin 1998, 2005; Rupke et al. 2005; Chisholm et al. 2015) that trace gas at \( T \sim 10^{4.5} \text{ K} \), but there has also been direct detection of kinematically confirmed hot winds (\( T \sim 10^{6.5} \text{ K} \)) via diffuse X-rays (e.g., Read et al. 1997; Lehmer et al. 1999; Strickland & Heckman 2007), as well as cold atomic and molecular outflows (e.g., Sturm et al. 2011; Bolatto et al. 2013; Leroy et al. 2015; Martini et al. 2018). Due to the low density and hence low emissivity of the hot gas, quantitative characterization of full multiphase outflows from observations has been limited to a few best-case examples (e.g., Strickland & Heckman 2007; Leroy et al. 2015); significant advances will require next-generation X-ray observatories (e.g., AXIS, ATHENA, and Lynx).

To date, systematic theoretical studies of outflow properties for different thermal phases have also been limited. Utilizing parsec-resolution local kiloparsec patches of galactic disks, resolved multiphase ISM simulations with SN feedback (and additional feedback processes) have been conducted by several groups. However, due to the complexity and expense of modeling full star-forming ISM physics with high resolution, many previous simulations studying galactic winds have adopted prescribed SN rates and positions (e.g., Creasey et al. 2013; Girichidis et al. 2016a, 2016b, 2018a; Martizzi et al. 2016; Li et al. 2017) and run only for a short period of time, with a limited range of ISM conditions (e.g., Gatto et al. 2017; Kannan et al. 2020). To understand galactic outflows as emergent phenomena produced by the star-forming ISM, lack of self-consistency is a concern because the reported characteristics could be sensitive to the adopted feedback rates and SN locations. Previous controlled simulations with SNe imposed “by hand” have shown that the resulting ISM and outflow properties change dramatically when SNe are located only in dense gas or randomly (e.g., Girichidis et al. 2016b), or when clustering of SNe is varied (e.g., Fielding et al. 2018). Simulations with a short duration are problematic because results may be strongly affected by imposed ISM initial conditions and numerical startup transients.

A different approach from high-resolution “local patch” simulations is global isolated galaxy and cosmological zoom-in simulations. In the case of cosmological zooms (e.g., Muratov et al. 2015; Christensen et al. 2016; Anglés-Alcázar et al. 2017; Tollet et al. 2019), cosmic accretion and merging/interaction of galaxies are included, which provides a “natural” CGM environment with which winds may interact (Fielding et al. 2020b). For studying wind acceleration, global/zoom-in models also have an advantage over local models in that the
effect of global geometry and quasi-conical wind expansion and acceleration is naturally captured (e.g., Chevalier & Clegg 1985). However, for studying wind creation, zoom-in simulations are at a disadvantage compared to local models in that the ISM physics, including star formation and feedback, is at best only marginally resolved. The adopted mass resolution ($\sim 10^{3-5} M_\odot$) is still insufficient to resolve the Sedov–Taylor stage of SNR evolution; since the remnant mass at the time of shell formation ($\sim 10^3 M_\odot$) must be resolved by several elements, mass resolution of $\lesssim 100 M_\odot$ is needed (Kim & Ostriker 2015a). This is critical for accurately modeling hot gas production and the multiphase interactions inherent to wind launching (Hu 2019). The derived wind properties from zoom-in simulations are compromised by approximate treatments of SN feedback; artificially delayed cooling (Christensen et al. 2016; Tollet et al. 2019) or momentum feedback (Muratov et al. 2015, 2017; Anglés-Alcázar et al. 2017).

Even in isolated galaxy simulations, achieving high enough resolution for resolving individual SNe, as well as self-consistent modeling of star formation with self-gravity, is challenging; currently such simulations are done only for very low mass galaxies (total gas mass $\sim 10^6 M_\odot$; e.g., Emerick et al. 2018; Hu 2019). For more massive galaxies, prescribed rates and positions of SNe are still adopted (e.g., Fielding et al. 2017; Schneider & Robertson 2018). Note that in their study of superwinds, Schneider et al. (2020) consider an isolated global galaxy with comparable resolution to local simulations (up to 5 pc) but with a much smaller set of physics: neither self-gravity nor magnetic fields, no cold ISM (cooling is truncated at $10^3$ K), and prescribed SN feedback rates and positions. Nevertheless, these simulations demonstrate the importance of uniformly high resolution in the extraplanar region for following both the hot and cool components of outflows and the interaction between phases.

Another potential issue in characterizing multiphase outflow properties from cosmological zoom-in or isolated global simulations is the adaptive resolution (either AMR or semi-Lagrangian method) that is usually employed (an exception is the CGOLS suite by Schneider & Robertson 2018). Although semi-Lagrangian/adaptive resolution (generally at constant mass) provides better resolution at higher densities to improve treatment of star formation, the low-density hot gas, which carries the majority of outflowing energy and metals, can be quite underresolved. The phase structure and overall energetics of the ISM and CGM depend sensitively on accurately resolving the mixing at interfaces between hot and cool gas (e.g., Fielding et al. 2020a), which would require extremely high mass resolution given the low density of the hot gas. Indeed, recent work employing fixed “spatial” rather than “mass” resolution in the CGM has revealed dramatic differences in multiphase gas properties (e.g., Hummels et al. 2019; Peebles et al. 2019; van de Voort et al. 2019). For winds, underresolution of the hot gas raises potential concerns about numerical phase mixing that could lead to underestimated metal loading or overestimated mass loading in an artificially phase-mixed wind, depending on the halo potential (see discussion in Kim & Ostriker 2018).

In the large-box cosmological simulations that are necessary for predicting statistics of galactic populations and for connecting baryonic distributions on large scales to cosmological parameters, typical mass resolutions are $> 10^5 M_\odot$ (see Figure 1 in Nelson et al. 2019 for a recent compilation). Even with significant improvements in computing power, it will continue to be necessary into the future to apply subgrid methodology in treating star formation and winds, because directly representing the physical processes involved would require several orders of magnitude higher resolution. Key wind parameterizations that are usually required by large-box cosmological simulations (as well as semianalytic cosmological models) are (1) the dimensionless mass loading, consisting of the mass (hydrogen and metals) carried out by the wind per stellar mass formed, and (2) the energy loading, consisting of the fraction of the original SN energy that is transferred to the outflowing gas. In addition, it is necessary to set (3) the wind velocity; this is often scaled relative to the halo velocity, but more generally a momentum loading (momentum ejection per stellar mass formed) or characteristic outflow velocity (or its distribution) can be given. Currently, the standard practice (e.g., Oppenheimer & Davé 2006; Vogelsberger et al. 2013; Crain et al. 2015; Pillepich et al. 2018) is to tune the wind parameters so that the resulting global galaxy properties match empirical constraints. Cosmological SAMs also adopt empirical prescriptions for the impact of stellar-driven winds on the ISM and CGM. Although SAM feedback prescriptions have traditionally been tuned to match observations, it is also interesting to study how different these prescriptions are from those that emerge from detailed numerical simulations (Pandya et al. 2020).

In the present paper, as part of the first results from SMAUG, we take a step toward the goal of a new subgrid approach by providing a detailed characterization of the mass, momentum, energy, and metal loading of multiphase outflows, based on parsec-resolution simulations of the star-forming ISM. In contrast to earlier local simulations, the major advance of our work is to achieve self-consistency in the ISM evolution, star formation, and feedback, as well as uniformly high resolution and long-term evolution. Using a new numerical framework called TIGRESS (Kim & Ostriker 2017, hereafter KO17), we resolve the self-gravitating collapse of star-forming cloud complexes, the energy-conserving stage of SNR evolution when hot gas is created, and the subsequent interactions between diffuse hot and denser warm gas. KO17 delineated the numerical methods involved and demonstrated their application by running a solar-neighborhood model over three orbit times, covering $\sim 10$ star formation/feedback cycles. Over this time, the ISM achieves a quasi-steady state with self-regulated star formation, insensitive to the initial setups. KO17 presented a thorough resolution study and confirmed convergence of turbulence amplitudes, thermal phase balance, magnetic field strength, SFRs, and outflow rates. Kim & Ostriker (2018, hereafter KO18) analyzed multiphase wind properties in solar-neighborhood TIGRESS simulations, focusing on the dichotomy of warm fountains and hot winds, loading factors as a function of heights and phases, and distributions of outflow velocities in the warm outflow.

The current paper focuses on the systematic investigation of outflowing gas, separately characterizing multiple phases of gas in a suite of seven TIGRESS simulations. We quantify mass, momentum, energy, and metal loading factors, as well as characteristic outflow velocities and metal enrichment factors. We follow KO17’s definition of thermal phases but merge the three lowest-temperature phases to a single cool

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8. https://www.simonsfoundation.org//flatiron/center-for-computational-astrophysics/galaxy-formation/smaug/papersplash1
(T < 2 × 10⁴ K) component, while keeping the intermediate (2 × 10³ K < T < 5 × 10⁵ K) and hot (T > 5 × 10⁶ K) phases. Physical conditions in our set of seven models span two orders of magnitude in gas surface density (Σgas ~ 1–100 M⊙pc⁻²) and four orders of magnitude in SFR surface density (ΣSFR ~ 10⁻⁴–1 M⊙ pc⁻² yr⁻¹), covering typical local conditions of nearby star-forming galaxies (e.g., Bolatto et al. 2017; see also Motwani et al. 2020, for our parameter space coverage in comparison with Illustris-TNG galaxies).

To further characterize outflows, in a companion paper we shall present joint probability distribution functions (PDFs) of gas outflow velocities and sound speeds. We shall also provide a guide to combining the loading factors presented here and PDFs and in order to apply them in large-scale models of galaxy formation.

The structure of this paper is as follows. Section 2 summarizes the TIGRESS numerical framework and introduces the model suite parameters employed in this paper. Section 3 explicates the cycles of star formation, feedback, and outflow/inflow that emerge in all the simulations. Section 4 presents multiphase outflow properties, including outflow fluxes, metal properties, and characteristic velocities, as well as average loading factors. Section 5 compares outflow characteristics with galactic properties, including SFR surface density (ΣSFR), gas surface density (Σgas), midplane density (ρmid), midplane pressure (Pmid), gas weight (W), and gas depletion time (τdep) to derive scaling relations. Section 6 discusses our results in the context of other existing simulations and observations, provides physical interpretations, and summarizes strengths and weaknesses of our numerical model. In Section 7, we summarize our results.

2. Methods and Models

This paper investigates properties of a suite of local ISM simulations in star-forming galactic disks to provide a comprehensive characterization of multiphase galactic outflows driven by stellar feedback. To evolve the ISM with star formation and stellar feedback self-consistently, we utilize the TIGRESS framework described in KO17. We refer the reader to KO17 for details of implementations and tests. In Section 2.1, we summarize key features and modifications of the TIGRESS framework from KO17. In Section 2.2, we introduce model parameters for our suite of simulations.

2.1. Methods

The TIGRESS framework evolves the ISM by solving the ideal MHD equations, including gravity and cooling/heating, in a local, rotating frame with a galactic orbital frequency Ω(R0) at a galactocentric distance R0. Local Cartesian coordinates x and y represent the local radial (R − R0) and azimuthal (R0φ − Ωt) directions, respectively, while z represents the vertical distance from the midplane. Shearing-periodic and outflow boundary conditions are adopted in the horizontal and vertical directions, respectively. We use the Athena finite-volume code for MHD (Stone et al. 2008; Stone & Gardiner 2009) with additional physics modules. The shearing box approach (Stone & Gardiner 2010) allows us to model the ISM in the context of rotating disk galaxies with uniformly high resolution (∼ O(1) pc) everywhere.

To follow star formation by gravitational collapse, the TIGRESS framework includes self-gravity by solving Poisson’s equation using fast Fourier transforms (Gammie 2001; Koyama & Ostriker 2009) and forms sink particles to represent star cluster formation in cells undergoing unresolved gravitational collapse (Gong & Ostriker 2013). The sink particles then further accrete if gas flows are converging into a virtual control volume (3³ cells surrounding a particle) from all three directions. Gas accretion onto a given sink particle ceases as soon as the first SN explodes (the SN event is stochastically determined, with the first event typically 3–4 Myr after the birth; see below). When sink particles are first formed or actively accreting, we reset the gas density, momentum, and pressure within the control volume with the extrapolated values from the nearby cells and dump only the difference between original and extrapolated values of mass and momentum in the control volume into the star particle. In the original Gong & Ostriker (2013) treatment adopted in KO17, all of the mass flux into the control volume is added to sink particles and the control volume is treated as ghost zones. Since control volume cells become active zones if a sink particle becomes a nonaccreting passive particle or merges with other particles, this approach is not strictly mass conservative. As initially applied in KO17, this nonconservation has a minimal effect in the total mass (net difference ~10% over 3tgas) of the R8 model, because the SFR is low and particle merging is not frequent, but it can be more significant for models with high SFRs. In the new approach, where the sink particle control volumes are treated as potential active cells, mass conservation is improved (see also Lam et al. 2019); for example, the cumulative effect in mass is at 3% over torb for model R4. By comparison, the total ISM mass reduction over 0.5 < t/torb < 1.5 is 23%, with 17% going into star formation and 8% into winds. Nonconservation is smaller for models with smaller SFRs and ~4%–5% for models R2 and LGR2. It should be kept in mind that instantaneous relationships among SFRs, outflow properties, and ISM properties are not affected by this slow secular variation, and the nonconservation does not affect any of the measures we report. In particular, all measures of outflows are obtained directly from fluxes in the simulation. The nonconservation of mass (reflecting a small addition from “reactivated” control volume cells) simply makes the mean value of Σgas at most a few percent larger than it would otherwise be over the simulation duration.

Stellar feedback in the TIGRESS framework includes the effects of far-UV (FUV) radiation and SN explosions. We slightly update the treatment of the heating rate due to FUV radiation from young stars. FUV radiation absorbed by small grains (e.g., PAHs) produces photoelectrons that heat the gas (Bakes & Tielens 1994). This is believed to be the dominant heating process in the neutral (atomic) ISM (Wolff et al. 1995), where FUV is not shielded (at low column densities) and where dust is not destroyed. To first order, the heating rate is proportional to the mean FUV intensity. Allowing for background heating from the metagalactic UV (Sternberg et al. 2002), the heating rate is given by

\[ \Gamma = \Gamma_0 \left( \frac{J_{\text{FUV}}}{J_{\text{FUV},0}} + 0.0024 \right), \tag{1} \]

where we adopt as reference solar-neighborhood values a heating rate of \( \Gamma_0 = 2 \times 10^{-26} \text{ erg s}^{-1} \) (Koyama & Inutsuka 2002) and a mean FUV intensity of \( 4\pi J_{\text{FUV},0} = 2.7 \times 10^{-3} \text{ erg s}^{-1} \text{ cm}^{-2} \) (or \( G_0 = 1.7 \); Draine 1978). We note that \( \Gamma_0 \) is held fixed,
implying that in the present treatment we do not allow for variations in dust abundance or photoelectric efficiency (Bakes & Tielens 1994; Weingartner & Draine 2001).

In the TIGRESS framework, the total FUV luminosity is calculated by summing up FUV luminosity of individual star clusters

$$L_{\text{FUV}} = \sum_i \Psi_{\text{FUV}}(t_{\text{age,sp}}) M_{\text{sp}}$$  \hspace{1cm} (2)

using a tabulated time-dependent mass-to-luminosity ratio $\Psi_{\text{FUV}}$ (from STARBURST99 as in Leitherer et al. 1999; see Figure 1 of KO17), star cluster age $t_{\text{age,sp}}$ (mass-weighted average is taken when there is addition of mass from accretion and merging), and star cluster mass $M_{\text{sp}}$. As our star clusters have masses $\gtrsim 10^4 M_\odot$, we adopt a fully sampled Kroupa initial mass function (IMF; Kroupa 2001) in setting $\Psi_{\text{FUV}}$. In KO17, we calculated the mean FUV intensity as $4\pi J_{\text{FUV}} = \Sigma_{\text{FUV}} = L_{\text{FUV}}/(L_\odot L_\odot)$, assuming uniformly spread radiation over the horizontal area of $L_\odot L_\odot$. This is valid for solar-neighbourhood and outer disks (e.g., R8, R16, and LGR8) but may overestimate the interstellar radiation field in denser environments where attenuation is generally higher. To allow for attenuation in an average sense, here we set the mean intensity based on the plane-parallel solution of the equation of radiation transfer in a slab with a uniform source distribution,

$$4\pi J_{\text{FUV}} = \Sigma_{\text{FUV}} \frac{(1 - E_2(\tau_L/2))}{\tau_L},$$  \hspace{1cm} (3)

where $\tau_L = \kappa_{\text{FUV}} \Sigma_{\text{gas}}$ is the UV−optical depth perpendicular to the slab and $E_2$ is the second exponential integral.\footnote{In follow-up work applying the adaptive ray-tracing method of Kim et al. (2017b), we are further testing this approximation.} We adopt $\kappa_{\text{FUV}} = 10^3$ cm$^2$ g$^{-1}$. Note that our heating rate as a result is time-varying but uniform in space, modulo a tweak at high temperatures $T > 10^5$ K.

The SN treatment is unchanged from KO17. When an SN explodes, we first calculate the mean gas properties (total mass $M_{\text{SNR}}$ and mean density $n_{\text{amb}}$) for the cells whose cell-centered distances from the explosion center are smaller than $R_{\text{SNR}} = 3.2 x$. We inject both thermal and kinetic energy with a ratio consistent with the Sedov–Taylor stage (0.72:0.28) if $M_{\text{SNR}}/M_\odot < 1$, where $M_\odot = 1540 M_\odot / (n_{\text{amb}}/$ cm$^{-3})^{-0.33}$ is the shell formation mass at a given ambient medium density $n_{\text{amb}}$ (Kim & Ostriker 2015a). If the shell formation mass is unresolved (i.e., $M_{\text{SNR}}/M_\odot > 1$ within the feedback region), we instead inject the terminal momentum of SNR $P_{\text{SNR}} = 2.8 \times 10^5 M_\odot$ km s$^{-1} / (n_{\text{amb}}/$ cm$^{-3})^{-0.17}$ (Kim & Ostriker 2015a). We find that more than 90% of SNe are well resolved (i.e., $M_{\text{SNR}}/M_\odot < 0.1$) in the simulations presented here.

With each SN explosion, we eject massless test particles with 50% probability to represent a runaway originating from a binary OB star. The ejection velocity follows an exponential distribution with exp($-\gamma_{\text{run}}/50$ km s$^{-1}$) for $\gamma_{\text{run}} \in (20, 200)$ km s$^{-1}$ (Eldridge et al. 2011), and the direction is chosen isotropically. Each runaway moves under the total gravitational potential and explodes as an SN after a pre-assigned explosion time. The total SN rate from a star cluster, including its runaways, is consistent with the SN rate from

STARBURST99 (Leitherer et al. 1999). KO18 showed that the outflow properties are not sensitive to the inclusion of runaways.\footnote{Andersson et al. (2020) explored the effect of runaways in isolated galaxy simulations and found large enhancement of mass outflow rates and corresponding loading factors ($\times 5-10$) when runaways were included. However, this result may reflect a numerical rather than a physical effect. In particular, it is possible that their no-runaways simulation failed to drive strong outflows because the majority of SNe had numerically unresolved evolution that failed to create hot gas that breaks out from the disk. The adopted AMR scheme (the RAMSES code) has a refinement strategy of splitting cells when the cell mass exceeds a designated maximum mass, in this case $\sim 4 \times 10^5 M_\odot$. The majority of SNe within star-forming, dense gas at the maximum level of refinement would then occur in cells with mass exceeding the shell formation mass $\sim 10^5 M_\odot$, with feedback implemented via momentum injection. However, runaway particles that have moved far from their birthplaces may be in lower-density environments, with higher mass refinement, at the time of SN explosions. To some extent, the inclusion of runaways is a partial solution to numerical difficulties in resolving SNe and driving hot superbubble breakout in moderate resolution simulations like Andersson et al. (2020). In our simulations, however, resolution is much higher and the majority of SNe in clusters (>90%) resolve the Sedov–Taylor stage of evolution, so that inclusion of runaways has insignificant impact on outflows.}

We use a tabulated cooling function from a combination of Koyama & Inutsuka (2002) at $T < 10^{4.2}$ K and Sutherland & Dopita (1993) at $T > 10^{4.2}$ K for solar metallicity. Although we include a metal tracer field, which is initialized with solar metallicity $Z_{\text{ISM,0}} = 0.02$, our cooling function does not depend on gas metallicity. In the model with the highest $Z_{\text{ISM}}$ (R2; see below), we find that the ISM metallicity reaches $\sim 0.04$ at the end of simulation, which is well within a range of observational estimates of the ISM metallicity in nearby star-forming galaxies (e.g., Lian et al. 2018). The hot outflow is typically more enriched by a factor of 2 than the ISM (although the cooling rate of the diffuse hot gas is quite low in any case). To address the effect of realistic, metallicity-dependent cooling, development of a second-generation TIGRESS framework, including more complete treatments of cooling, radiation, and chemistry, is now underway.

### 2.2. Models

For this paper, we use TIGRESS runs with seven different parameter sets, covering conditions generally representative of inner and outer regions of Milky Way−like galaxies, including the solar-neighbourhood model described in KO17 and KO18. We list key parameters of these models in Table 1. The gas surface density $\Sigma_{\text{gas,0}}$ in Column (2) is the initial value in the simulation and decreases over time because gas turns into sink particles owing to star formation and escapes vertically as a wind. SNe return mass to the gas in the form of ejecta, but on average this is only 10% of that locked into stars. The galactic environment parameters, such as angular speed of galactic rotation $\Omega$, stellar surface density $\Sigma_{*}$, stellar scale height $z_*$, dark matter halo density $\rho_{\text{dm}}$, and galactocentric radius $R_0$, are fixed in time for each simulation; these parameters are
important for setting the gravitational potential (see KO17 for the analytic expression) and for setting the differential shear rate (important to dynamo activity; e.g., Käpylä et al. 2018) and Coriolis force. The “LG” models have external (stellar and dark matter) vertical gravity reduced by about a factor of 8 near the midplane \((z < \zrc)\) and 4–5 far above the disk \((z > \zrc)\) compared to the corresponding R2, R4, and R8 models. All simulations use \((N_x, N_y, N_z) = (256, 256, 1792)\) zones and uniform cubic grid cells with side length of \(\Delta x\) (Column (8)). Our parameter choice covers typical ranges seen in nearby star-forming galaxies (e.g., Sun et al. 2020), as well as cosmological simulations (Motwani et al. 2020).

Additional parameters are used to set initial conditions of the gas in the simulation, including the temperature profile and turbulent vertical velocity dispersion \(\sigma_{z,0}\) and plasma beta \(\beta_0 \equiv P/P_{\text{mag}}\) (see KO17). The initial vertical profiles of density, pressure, and azimuthal magnetic field are set to be in a hydrostatic equilibrium with given total vertical velocity dispersion, including thermal, turbulent, and magnetic terms. However, the initial thermal and turbulent support is lost quickly owing to radiative cooling and turbulence dissipation. The gas soon falls toward the midplane, and this density increase triggers a burst of star formation. As we shall discuss in Section 3, the first burst is not fully self-consistent because it is subject to the initialization. Over time, the evolution becomes self-regulated; our analysis therefore will focus on the time subsequent to the first burst. To offset the rapid initial cooling and turbulence dissipation, we introduce randomly placed star particles in the initial conditions with age and mass distributions corresponding to the \(\Sigma_{\text{SFR}}\) at later times (estimated from lower-resolution simulations). We adopt \(\beta_0 = 10\) for all models, except R2 with \(\beta_0 = 2\), and \(\sigma_{z,0} = 30, 15, 10\), and 10 km s\(^{-1}\) for R2/LGR2, R4/LGR4, R8/LGR8, and R16.

From several independent simulation runs with different \(\sigma_{z,0}\) and \(\beta_0\) (typical ranges are \(\sigma_{z,0} = 10–30\) km s\(^{-1}\) and \(\beta_0 = 1–10\)), we have confirmed that the evolution is statistically converged irrespective of initial conditions unless the initial parameters are extreme; the initial magnetic field strength can impact the overall outcomes if it is too strong or too weak compared to the saturated value since the evolution of the regular magnetic field is much slower than all other timescales (see Kim & Ostriker 2015b).

Magnetic fields in outflows do not contribute to momentum and energy fluxes significantly. However, inclusion of magnetic fields has indirect effects on outflows and associated galactic properties by increasing the vertical pressure support near the midplane and reducing SFRs. We find that contribution from magnetic stresses to the vertical pressure support can be as high as 50% (depending on initial field strengths since saturation is not achieved within \(1–2\) \(t_{\text{orb}}\)), but typically about 30% (E. Ostriker & C.-G. Kim 2020, in preparation).

The numerical in each model name indicates the galactocentric radius of the simulation box; e.g., the box in model R8 is centered at \(R_0 = 8\) kpc. The spatial resolution in parsecs is progressively smaller from model R16 to model R2 (also implying a smaller simulation box) as we move from outer (lower-density) to inner (higher-density) galactic regions. At higher densities, both thermal and dynamical length scales are smaller. For each model, we tested varying simulation box sizes, and the values ultimately adopted were optimized such that resolution is sufficiently high while still providing a large enough horizontal area such that superbubbles do not fill the entire horizontal domain. For our standard simulations, the horizontal box size, \(L_x\) and \(L_y\), decreases from 2048 pc for model R16, to 1024 pc for model R8, to 512 pc for models R2 and R4. In Appendix A, we briefly discuss the role of box size and show the resolution dependence of our results to demonstrate convergence.

Finally, we note that although a value of \(R_0\) is adopted for each model, this is only used in setting the local background rotational velocity and the shape of dark matter halo gravity; the simulations are all local, and in principle could equally well describe similar conditions within a dwarf as a massive spiral (at a given metallicity).

3. Overall Evolution

In previous papers (KO17 and KO18), we have presented the overall evolution of the solar-neighborhood model (R8 in this paper). The evolution exhibits multiple feedback cycles, reaching a quasi-steady state in which the SFR is self-regulated by stellar feedback. We will present a comprehensive analysis of the suite of simulations presented here in context of the theory of pressure-regulated, feedback-modulated star formation in a separate paper (E. Ostriker & C.-G. Kim 2020, in preparation; see also Ostriker et al. 2010; Kim et al. 2011; Ostriker & Shetty 2011). In this paper, our focus is mainly on outflows from the main gas layer, above the scale height of gas. However, here we also briefly cover star formation self-regulation since the bursts and lulls of star formation are responsible for the cyclic behavior of outflow/inflow.

As soon as the simulation begins, the initial turbulent energy begins to dissipate and the denser gas cools to form the cold medium. Material falls vertically, and dense, cold cloud complexes form and collect near the midplane. Star clusters are born in gravitationally collapsing parts of the cloud complexes; these heat the ISM by emitting UV radiation and drive turbulence through SN explosions, restoring the lost vertical support. As the disk puffs up, the overall SFR drops. The now-reduced stellar feedback cannot offset cooling and turbulent dissipation, so that gas falls back to the midplane and the next star formation event follows. The cycle repeats, with the system entering a self-regulated, quasi-steady state.

In this state, the time-averaged total vertical pressure support (sustained by star formation feedback) balances the vertical weight of gas (as shown in simulations of Kim et al. 2013; Kim & Ostriker 2015b; see also Thompson et al. 2005;
However, instantaneously there is always a mismatch between “supply” (pressure from feedback) and “demand” (weight from gravity). Especially, the injection of energy and momentum from SN feedback is highly concentrated in space and time, leading to an overshoot (e.g., Benincasa et al. 2016; Orr et al. 2019). The resulting gas outflows carry the excess momentum and energy into the extraplanar region (above the gas scale height), and a portion is eventually vented to the CGM. As we shall show (see also KO18), in our simulation suite the hot gas created by SN shocks is mainly responsible for energy and momentum delivery to the extraplanar region and beyond, while the cooler gas delivers significant mass beyond the disk scale height. The self-regulation and outflow cycle is evident in all models, with some qualitative differences.

To help visualize feedback-driven outflows in the simulation suite, Figures 1 and 2 show density and temperature slices, respectively, at \( y = 0 \) for snapshots at \( t/t_{\text{orb}} = 0.7 \). Also, in Figure 1 we show velocity streamlines color-coded by outward vertical velocity,

\[
\mathbf{v}_{\text{out}} \equiv v_z \text{sgn}(z).
\]

The hot, fast outflows preferentially vent through low-density chimneys, carved out of the denser warm ISM by superbubble breakout events. At the same time, a highly dynamic fountain...
of clumpy, cooler gas coexists with hot gas in the extraplanar region and is both inflowing and outflowing. Turbulent flows of cool gas close off chimneys, limiting hot outflows and leading to significant interaction between hot winds and cool fountains.

Figure 3 shows the horizontally averaged vertical mass flux, $\langle \rho v_{\text{out}} \rangle_{xy}$, from all models as a time series. The spacetime diagram of mass flux profiles demonstrates the cycles of outflow/inflow. Outward fluxes ($\langle \rho v_{\text{out}} \rangle_{x,y} > 0$) are in red, while inward fluxes ($\langle \rho v_{\text{out}} \rangle_{x,y} < 0$) are in blue. We separate gas using three temperature bins, $T < 2 \times 10^4$ K for cool (left column), $2 \times 10^4 \text{K} < T < 5 \times 10^5$ K for intermediate (middle column), and $T > 5 \times 10^5$ K for hot gas (right column). For reference, we plot $H$ and $2H$ as solid and dashed lines, respectively, in the middle column, where $H$ is the instantaneous scale height of gas (the mass-weighted dispersion of vertical gas positions; see Equation (8)).

Focusing on the left column, where we show the “cool” component, cyclic behavior of alternating outflow and inflow is evident for all models. The evolution is more regular for R16 and gets more complex at higher surface densities. For R8, LGR4, and LGR8, the evolution is still quite cyclic, while R2, R4, and LGR2 show complex interaction between outflows and inflows and generally less cyclic evolution, especially for the gas near the midplane.

Qualitative differences in the cyclic behavior among models can be understood from competition between key timescales: the vertical oscillation time $t_{\text{osc}}$ and the star cluster evolution timescale $t_{\text{evol}} \sim 40$ Myr (Leitherer et al. 1999). The former

---

Figure 2. Same as Figure 1, but for temperature slices. The video begins at $t/t_{\text{orb}} = 0$ and ends at $t/t_{\text{orb}} = 1.48$. The real-time duration of the video is 20 s. It can also be found at changgo.github.io/tigrass-wind-figureset/movies.html. (An animation of this figure is available.)

---

11 The cool phase in this paper includes cold/unstable/warm gas (or two-phase gas) as defined in KO17 and KO18. Although we do not explicitly distinguish cold, unstable, and warm gas as in previous work, the fractions of cold and unstable components are negligible at $|z| > H$. Therefore, the cool phase in this paper is essentially equivalent to the warm gas of KO18.
controls the self-regulation cycle because gas pushed outward by feedback from a burst of star formation returns after \( t_{\text{osc}} \) and participates in the next star formation event. The latter sets the duration of energy/momentum injection from a given star formation event, during which star formation is generally reduced.

In Table 2, we list three timescales along with relevant quantities measured from simulations to obtain these timescales. Column (2) gives the orbital period of galactic rotation \( t_{\text{orb}} \) with the usual definition:

\[
t_{\text{orb}} \equiv \frac{2\pi}{\Omega} = 120 \text{ Myr} \left( \frac{\Omega}{50 \text{ km s}^{-1} \text{kpc}^{-1}} \right)^{-1}.
\]

For \( t_{\text{osc}} \), we list both a measure from the simulation and an analytic estimate in Columns (3) and (4), respectively. The
vertical gravity is nearly linear $g_z \approx -4\pi G \rho_{\text{gas}} \langle z \rangle$ for the majority of gas since the gas scale height is smaller than or comparable to the stellar height ($z_\ast$) and dark matter scale length ($R_0$) assumed in our potential model (except for model R16). The collisionless vertical oscillation time can then be approximated only in terms of the total midplane density $\rho_{\text{tot}} = \rho_{\text{gas}} + \Sigma_*/(2z_\ast) + \rho_{\text{dm}}$ as

$$t_{\text{osc},a} \approx \frac{2\pi}{(4\pi G \rho_{\text{tot}})^{1/2}} = 37 \text{ Myr} \left(\frac{\rho_{\text{tot}}}{0.5 M_\odot \text{ pc}^{-3}}\right)^{-1/2}. \quad (6)$$

Note that the midplane density of the gas $\rho_{\text{gas}}$ (or $n_{\text{gas}} = \rho_{\text{gas}}/(\mu m_H)$) is calculated by taking the mean of density in the two horizontal planes at $z = \pm \Delta z_\ast/2$. Since the prediction for the scale height under linear gravity is $H = \sigma_{\epsilon}/(4\pi G \rho_{\text{tot}})^{1/2}$, an alternative definition of the vertical oscillation time measurable directly from gas properties is

$$t_{\text{osc},n} \equiv \frac{2\pi H}{\sigma_{\epsilon,\text{eff}}} = 46 \text{ Myr} \left(\frac{H}{300 \text{ pc}}\right) \left(\frac{\sigma_{\epsilon,\text{eff}}}{40 \text{ km s}^{-1}}\right)^{-1}. \quad (7)$$

Here we calculate $H$ and $\sigma_{\epsilon,\text{eff}}$ from the mass-weighted height and effective velocity dispersion measured in the simulation, where they are respectively defined by

$$H \equiv \left(\int \rho v^2 dV\right)^{1/2} / \int \rho dV \quad (8)$$

and

$$\sigma_{\epsilon,\text{eff}} \equiv \left(\frac{\int [\rho v_x^2 + P + B_z^2/(8\pi) - B_z^2/(4\pi)] dV}{\int \rho dV}\right)^{1/2}. \quad (9)$$

Note that the effective vertical velocity dispersion includes the contributions from turbulent, thermal, and magnetic stresses. The values are all time averages over $0.5t_{\text{orb}} < t < 1.5t_{\text{orb}}$. $t_{\text{osc},n} \sim 0.5t_{\text{orb}}$ for our models, except R16.

If $t_{\text{osc}}$ is sufficiently longer than $t_{\text{evol}}$ (e.g., for R16), a major star formation event cannot occur until after previously blown-out gas falls back. If $t_{\text{osc}}$ is smaller than or comparable to $t_{\text{evol}}$ (e.g., R2, R4, and LGR2), the situation is more complicated. Since each major star formation event continuously injects energy/momentum for $t_{\text{evol}}$, SN rates do not decline sharply for $t_{\text{evol}}$, gas that is launched and returns after $t_{\text{osc}}$ can be relaunched before participating in the next star formation event. The self-regulation cycle is delayed until feedback shuts off after $\sim t_{\text{evol}}$.

The distinct oscillatory behavior seen in Figure 3 is in part due to the limited horizontal domain of the TIGRESS simulations. Because the natural horizontal correlation scale of star formation is not extremely small compared to the size of our simulation domain, averages at a given time will not statistically sample many independent regions at different stages of the evolutionary cycle. Synchronization within a local patch can also be enhanced if initial conditions tend to trigger a collapse of the entire disk, as in models LGR2 and LGR4. For LGR4, where $t_{\text{osc}} > t_{\text{evol}}$, the prominent oscillation cycle persists for a long time. However, for LGR2, even when the initial collapse induces very coherent first outflows, the feedback regulation cycles become highly irregular since $t_{\text{osc}} \sim t_{\text{evol}}$, so that the infalling gas keeps interacting with outflows from previous feedback events. On the other hand, while the early evolution of LGR8 is quite irregular, it eventually shows a fairly regular oscillation at later times since $t_{\text{osc}} \gg t_{\text{evol}}$. Overall, the late-time evolution ($t > 0.5t_{\text{orb}}$) and regularity of the cyclic behavior are self-consistently set by the fundamental timescales of the system.

The orbital time of galactic rotation $t_{\text{orb}}$ is relevant to the growth of the large-scale gravitational instability, due to the effects of epicyclic oscillations and shear (e.g., Goldreich & Lynden-Bell 1965; Elmegreen 1987; Kim & Ostriker 2001; Kim et al. 2002). In general, the gravitational timescale $t_g \sim \sigma/[G \Sigma]$ must be shorter than the epicyclic or shear times ($\sim t_{\text{orb}}$) for gravitational instabilities to grow. Typically, $t_g \sim t_{\text{orb}}$ in normal galaxies, i.e., the Toomre parameter is of order unity (Toomre 1964). If gravitational instability were the only important dynamical process acting on large scales, the inevitable result would be a strong starburst. However, in our simulations, $t_{\text{orb}}$ does not control star formation by itself because $t_{\text{osc}} < t_{\text{evol}}$, so that coherent structures at large scales are not able to continue growing for very long periods. Instead, they are destroyed by feedback before high star formation efficiency is achieved. We note, however, that in model R2 $t_{\text{orb}}$, $t_g \sim t_{\text{evol}}$, so that feedback is less able to limit large-scale gravitational instability. In reality, conditions with

### Table 2: Timescales and Relevant Measured Quantities

| Model | $t_{\text{tot}}$ (Myr) | $t_{\text{osc},a}$ (Myr) | $t_{\text{osc},n}$ (Myr) | $t_{\text{evol}}$ (Gyr) | $H$ (pc) | $\sigma_{\epsilon,\text{eff}}$ (km s$^{-1}$) | $n_{\text{gas}}$ (cm$^{-3}$) | $\rho_{\text{tot}}$ ($M_\odot$ pc$^{-3}$) | $\Sigma_{\text{gas}}$ ($M_\odot$ pc$^{-2}$) | $\Sigma_{\text{surf}}$ ($M_\odot$ kpc$^{-2}$ yr$^{-1}$) |
|-------|-----------------|-----------------|-----------------|-----------------|---------|-----------------|-----------------|-----------------|-----------------|-----------------|
| R2    | 61              | 32              | 23              | 6.6 $\times$ 10$^{-2}$ | 64      | 7.7             | 1.3             | 74              | 1.1             |                 |
| R4    | $1.1 \times 10^2$ | 51              | 37              | 0.23            | 41      | 1.4             | 0.50            | 30              | 0.13            |                 |
| R8    | 2.2 $\times$ 10$^2$ | 51 $\times$ 10$^2$ | 3.1 $\times$ 10$^2$ | 3.1 $\times$ 10$^2$ | 18 $\times$ 10$^2$ | 0.86            | 0.12            | 11              | 5.1 $\times$ 10$^{-3}$ |                 |
| R16   | 5.2 $\times$ 10$^2$ | 4.5 $\times$ 10$^2$ | 3.1 $\times$ 10$^2$ | 3.1 $\times$ 10$^2$ | 11 $\times$ 10$^2$ | 0.86            | 0.12            | 11              | 5.1 $\times$ 10$^{-3}$ |                 |
| LGR2  | 1.2 $\times$ 10$^2$ | 48              | 0.15            | 3.6 $\times$ 10$^2$ | 43      | 5.1             | 0.31            | 0.15            | 9.0 $\times$ 10$^{-2}$ |                 |
| LGR4  | 2.0 $\times$ 10$^2$ | 80              | 4.2 $\times$ 10$^2$ | 4.2 $\times$ 10$^2$ | 30      | 1.5             | 0.11            | 0.38            | 9.0 $\times$ 10$^{-2}$ |                 |
| LGR8  | 4.1 $\times$ 10$^2$ | 2.2 $\times$ 10$^2$ | 1.7 $\times$ 10$^2$ | 1.7 $\times$ 10$^2$ | 17      | 0.37            | 2.5 $\times$ 10$^{-2}$ | 10              | 3.2 $\times$ 10$^{-3}$ |                 |

Note: Column (2): orbit time (Equation (5)). Column (3): vertical oscillation time defined by Equation (7) using numerically measured gas scale height (Column (6)) and velocity dispersion (Column (7)). Column (4): vertical oscillation time defined by Equation (6) using total mass density at the midplane (Column (9)). Column (5): gas depletion time defined by the ratio of gas surface density (Column (10)) and SFR surface density (Column (11)). Column (6): gas scale height (Equation (8)). Column (7): effective vertical velocity dispersion (Equation (9)). Column (8): midplane number density of gas. Column (9): total midplane mass density of gas, stars, and dark matter. Column (10): gas surface density. Column (11): SFR surface density. Numerically measured quantities are averaged over $0.5t_{\text{orb}} < t < 1.5t_{\text{orb}}$. 

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very short orbital and gravitational timescales may also be subject to strong radial flows. Following this in detail would require global modeling, but unfortunately this is not yet tractable with the same uniformly high spatial resolution as our simulations.

Finally, we note that the gas depletion time is generally longer than \( t_{\text{scarc}} \leq t_{\text{vol}} \), and \( t_{\text{orb}} \), so that secular evolution has a minimal effect on the average properties in the self-regulated state.

In Figure 3, it is generally possible to link cool outflows (left) with the outflows of intermediate (middle) and hot (right) phases. Simultaneous, distinct outflows in all phases are realized when there is breakout of superbubbles produced by spatially and temporally correlated SNe. In an “outflowing” epoch, the hot outflows easily reach the domain boundaries without significant loss of mass flux. However, the cool gas launched with the hot gas after a burst eventually falls back. That is, red turns to blue for the cool gas. Notably, even during an “inflowing” epoch of the cool gas, the high-entropy hot gas continues to rise. That is, the hot gas shows only outflows (red), with no returning mass flux. We note, however, that even if there is a reasonably high SN rate producing hot gas, hot outflows are sometimes blocked by returning inflows of cool gas and cannot reach the boundary (see, e.g., R2 (top row) at \( t/t_{\text{orb}} \sim 1 \)). Thus, successful breakout is not solely determined by the SFRs (or SN rates) but is subject to the complex interaction between superbubble expansion and inflowing gas from previous events (sometimes in neighboring regions).

To summarize: outflows in our simulation suite show both regular and complex behaviors, depending on the model parameters (Figure 3). In the extrapolar region, outflows and inflows coexist in different phases, which we resolve in our simulations (Figures 1 and 2). In what follows, we explain how we characterize key properties of outflows and relate them to global properties, thereby deriving scaling relations.

### 4. Characterizations of Multiphase Outflows

In this section, we present characterization of multiphase outflows using outward mass, momentum, energy, and metal fluxes, separating the different thermal phases. We first present results for time evolution of outward fluxes (Section 4.1) and metal properties (Section 4.2) through surfaces (both upper and lower sides of the disk) at different heights, including two fixed heights at 500 pc and 1 kpc, and two time-dependent heights using the instantaneous gas scale height at \( H \) and \( 2H \).\(^{12}\)

We then show time-averaged vertical profiles of loading factors (Section 4.3), outflow velocities, and metal properties (Section 4.4). Figures and tables in this section are for model R4 or for values at \( |z| = H \). Figures for other models and the

\(^{12}\)The main motivation of this work is to report emergent multiphase outflow properties from resolved, self-consistent simulations of the star-forming ISM. In this undertaking, there is a tension between competing desiderata. On the one hand, it may be desired to measure outflowing gas properties at heights far from the disk midplane, where interactions with the “ISM” gas have been left behind. Larger distances are also closer to the resolution of big-box cosmological simulations. On the other hand, there is a countervailing need to choose a height closer to the midplane where the local approximation is valid (and climbing out of the global potential has not affected the outflow velocities). We shall show that interactions are minimized above \( \sim 2H \) while the local assumption is reasonable (with the local potential dominating over the global disk + halo potential and the flow streamlines not affected by global geometry) up to \( H \) or \( 2H \). Locations between \( H \) and \( 2H \) are thus a good compromise for making our measurements of outflow properties.

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data for tables at different heights are available at [https://changgoo.github.io/tigress-wind-figureset/figureset.html](https://changgoo.github.io/tigress-wind-figureset/figureset.html) and doi:10.5281/zenodo.3872049.

#### 4.1. Outflow Fluxes

The instantaneous outflow fluxes of each thermal phase through a horizontal surface at height \( z \) are calculated by summing up the vertical fluxes of cells with a positive outward vertical velocity \( v_{\text{out}} > 0 \) and temperature in the specified temperature range of each phase. Formally, the outflow flux of the quantity “\( \phi \)” in phase “\( \text{ph} \)” is defined by summing up the phase-selected flux over the horizontal domain with an area of \( L_xL_y \) as

\[
\mathcal{F}_{\phi,\text{ph}} = \sum_{i,j} F_{\phi}(i, j; t) \Theta(C) \frac{\Delta x \Delta y}{L_xL_y},
\]

where \( (i, j, k) \) is grid zone index, \( \Delta x = \Delta y \) is grid resolution (Column 8 in Table 1), and \( \Theta(C) \) is the top-hat-like filter that returns 1 if the conditional argument is true or 0 otherwise. Here the conditional argument is \( v_{\text{out}} \) > 0 for the outflowing gas and \( T \) in one of three temperature bins \( T < 2 \times 10^4 \text{ K} \), \( 2 \times 10^4 \text{ K} < T < 5 \times 10^4 \text{ K} \), and \( T > 5 \times 10^5 \text{ K} \) for the cool, intermediate, and hot phases, respectively. We consider four physical quantities \( q = M, p, E, \) and \( M_Z \) (we simply use \( Z \) in the subscript for succinctness; e.g., \( \mathcal{F}_Z \) is not metallicity flux but metal mass flux) to denote mass, \( z \)-momentum, energy, and metal mass, respectively. The corresponding vertical outgoing fluxes are defined by

\[
\mathcal{F}_M = \rho v_{\text{out}}, \tag{11}
\]

\[
\mathcal{F}_p = \rho v_{\text{out}}^2 + P, \tag{12}
\]

\[
\mathcal{F}_E = \rho (\frac{1}{2} v^2 + \frac{\gamma}{\gamma - 1} c_s^2), \tag{13}
\]

\[
\mathcal{F}_Z = \rho Z v_{\text{out}}, \tag{14}
\]

where \( v^2 = (v \cdot v) \), \( c_s^2 \equiv P/\rho \), and \( \gamma = 5/3 \), while other symbols have their usual meaning. We note that Equations (12) and (13) include both the kinetic and thermal components of the vertical momentum and total energy fluxes, ignoring the magnetic terms that are negligible in outflows. We analyze the total momentum flux instead of just the kinetic term separately since the contribution from the thermal pressure in the hot gas is substantial in outflows. The thermal pressure in cool gas is largely set by the balance between photoheating and cooling (rather than being driven by SNe) but is negligible in outflows. The simulation has nonzero metallicity in the beginning (\( Z_{\text{ISM},0} = 0.02 \)), and the metal mass flux consists of two origins, metals from the ISM and newly injected by SNe (\( Z_{\text{SN}} = 0.2 \)). Although we separately trace the total and SN-origin metals using independent passive scalar variables, separation between SN-origin and ISM-origin metals is nontrivial owing to enrichment of the ISM and recycling. We discuss this in more detail in the next section (Section 4.2). For now, we show the total metal flux.

Figure 4 plots, from top to bottom, the horizontally averaged mass, momentum, energy, and metal fluxes of model R4. We report combined fluxes through both upper and lower horizontal surfaces at a fixed height (dark and light blue
for $|z| = 0.5$ and $1$ kpc, respectively) or at a time-dependent height (dark and light red for $|z| = H$ and $2H$, respectively; see gray solid and dashed lines in the middle column of Figure 3 for the time variation of the scale heights). From the left to right column, we show cool, intermediate, and hot outflows separately.

From Figure 4, it is evident that the features of multiphase outflows at launch seen in the solar-neighborhood TIGRESS model (R8) (as reported by KO18) are generic across galactic conditions considered in our model suite. The cool component delivers most of the mass to the extraplanar region; in these simulations this cool gas subsequently returns to the midplane (the cool “fountain” is clear in the left column of Figure 3), but in a shallower global potential this cool outflow could escape.

The hot component carries most of the energy and escapes from the simulation domain as a wind. The intermediate component is subdominant in all fluxes. Due to the short cooling time of the intermediate component, a significant fraction of the outflow in this temperature range cools and mixes into the cooler gas in the course of its evolution (e.g., Vijayan et al. 2020, for more quantitative analyses). The fluxes, especially for the cool component, generally decrease for $|z| = 0.5$ and 1 kpc, respectively) or at a time-dependent height (dark and light red for $|z| = H$ and $2H$, respectively; see gray solid and dashed lines in the middle column of Figure 3 for the time variation of the scale heights). From the left to right column, we show cool, intermediate, and hot outflows separately.

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with distance from the midplane. Occasionally (e.g., at around \( t \sim 0.7\,\text{r}_{\text{SNe}} \); see also Figure 3), cool inflows are strong enough to shut off outflows nearly completely, showing dramatic drops from lower heights \(|z| = H\) or 500 pc; darker lines) to upper heights \(|z| = 2H\) or 1 kpc; lighter lines).

In every panel of Figure 4, the gray solid line shows a corresponding reference flux calculated based on the instantaneous SN rate, defined by

\[
\mathcal{F}_{\text{q,ref}} \equiv q_{\text{ref}} \frac{N_{\text{SN}}}{L_x L_y}. \tag{15}
\]

Here \(N_{\text{SN}}\) is the instantaneous SN rate calculated with an adaptive time window, within which the number of SNe is 100. The coefficient \(q_{\text{ref}}\) adopted in each reference flux is set based on simple physical considerations, as follows:

\[
M_{\text{ref}} = m_s = 95.5\,M_\odot \tag{16}
\]

\[
p_{\text{ref}} = E_{\text{SN}}/(2v_{\text{cool}}) = 1.25 \times 10^5\,M_\odot\,\text{km}\,\text{s}^{-1} \tag{17}
\]

\[
E_{\text{ref}} = E_{\text{SN}} = 10^{51}\text{erg} \tag{18}
\]

\[
M_{E,\text{ref}} = M_\odot Z_{\text{SN}} = 2\,M_\odot. \tag{19}
\]

Here we adopt a total mass of new stars formed for each SN of \(m_s = 95.5\,M_\odot\) (Kroupa 2001), an SN explosion energy of \(E_{\text{SN}} = 10^{51}\text{erg}\), a mean mass in ejecta from each SN of \(M_{\text{ej}} = 10\,M_\odot\), and a mean SN ejecta metallicity of \(Z_{\text{SN}} = 0.2\). We note that the combination of Equations (15) and (16) is equivalent to a reference mass flux of \(\mathcal{F}_{\text{M,ref}} = \Sigma_{\text{SFR}}\), the mean SFR per unit area averaged over the star cluster lifetime. Because SNe from a given star cluster persist over \(t_{\text{evol}} \sim 40\,\text{Myr}\) for a fully sampled IMF, the reference fluxes defined by Equation (15) depend on the SFR over the last 40 Myr and are therefore smoother than they would be if an instantaneous (or time-delayed) value of \(\Sigma_{\text{SFR}}\) were employed, while still giving the same long-term average.

For the reference momentum per SN, we adopt a value \(E_{\text{SN}}/v_{\text{cool}}\) with \(v_{\text{cool}} = 200\,\text{km}\,\text{s}^{-1}\), which represents the spherical momentum at the end of the Sedov stage when an SN blast wave cools and a shell forms (Draine 2011; Kim & Ostriker 2015a), also applying the geometric factor 1/2 to account for the vertical component of midplane-centered sources (Ostriker & Shetty 2011). An alternative reference momentum choice that is sometimes adopted is the initial SN ejecta momentum \(p_{\text{ej}} \equiv (2M_{\text{ej}}E_{\text{SN}})^{1/2} = 3.2 \times 10^4\,M_\odot\,\text{km}\,\text{s}^{-1}\). This (reduced by a factor 2) would be a more instructive choice if the SNR evolution remains in the free expansion stage until it reaches the height where a wind is launched. We generally find that this is not the case. We note that \(p_{\text{ref}}\) is an order of magnitude greater than the vertical momentum from initial SN ejecta, \(p_{\text{ej}}/2\).

From a large number of recent investigations of individual SNR evolution in inhomogeneous environments (e.g., Iffrig & Hennebelle 2015; Kim & Ostriker 2015a; Martizzi et al. 2015; Walch & Naab 2015) and of superbubble evolution driven by multiple SNe (e.g., Kim et al. 2017a; Fielding et al. 2018; El-Badry et al. 2019; Gentry et al. 2019), there is an emerging consensus in the community that the momentum injection to the ISM per event is relatively insensitive to ambient conditions. In particular, the momentum depends only very weakly on density, as shown in earlier uniform-background simulations (e.g., Cioffi et al. 1988; Blondin et al. 1998; Thornton et al. 1998). The terminal momentum per SN from clustered SNe is comparable to that from a single SN event as long as shocks from individual SNe remain supersonic, but it can be factor of a few smaller if blast waves from individual SNe become subsonic before reaching the shell (e.g., Kim et al. 2017a; El-Badry et al. 2019). Since most SNe are clustered, with the reference value of Equation (17) we expect the kinetic momentum loading near the midplane to be smaller than unity, and to be lower in models with higher SFR.

The dimensionless ratios between measured and reference fluxes are often termed loading factors (e.g., Somerville & Davé 2015; see Section 4.3). Although actual and reference fluxes in Figure 4 share similar evolutionary trends, the reference fluxes do not show the same large modulations as some of the measured fluxes. As we discussed in Section 3, complicated interaction between outflows and inflows makes one-to-one correspondence between strength of feedback (outflow driving) and emergent fluxes nontrivial (see also Appendix B).

4.2. Metallicity and Enrichment of Outflows

To understand the role of SN feedback in metal evolution within and beyond galaxies, simply measuring the total metal flux is insufficient. Every SN explosion injects metal mass \(Z_{\text{SN}}M_{\text{ej}} = 2\,M_\odot\), some of which goes directly to the extraplanar region as outflows, and some of which mixes with the ISM near the midplane, which was initialized with solar metallicity \(Z_{\text{ISM0}} = 0.02\). At a given epoch, outflowing gas can thus originate from one of three components: the ISM at the beginning of the simulation, \(\dot{M}(\text{ISM} \rightarrow \text{out})\), and SNe from previous SN events that have mixed into the ISM, \(\dot{M}(\text{SN} \rightarrow \text{ISM} \rightarrow \text{out})\), and SNe from \(\text{current}\) SN events, \(\dot{M}(\text{SN} \rightarrow \text{out})\). Note that \(\dot{M}(\text{SN} \rightarrow \text{ISM} \rightarrow \text{out})\) in principle includes metals recycled from fountain flows, which we do not separately track in this study. The total mass and metal outflow rates can be respectively written as

\[
\dot{M} = \dot{M}(\text{ISM} \rightarrow \text{out}) + \dot{M}(\text{SN} \rightarrow \text{ISM} \rightarrow \text{out}) + \dot{M}(\text{SN} \rightarrow \text{out}) \tag{20}
\]

\[
\dot{M}_Z = Z_{\text{ISM}}\dot{M}(\text{ISM} \rightarrow \text{out}) + Z_{\text{SN}}\dot{M}(\text{SN} \rightarrow \text{ISM} \rightarrow \text{out}) + Z_{\text{SN}}\dot{M}(\text{SN} \rightarrow \text{out}). \tag{21}
\]

In TIGRESS, we employ passive scalars for total and SN-injected metals, with densities that evolve under the mass conservation equation with a given velocity field. The total metal scalar allows us to measure \(M_{\text{Z}}\), while the SN-injected scalar traces the sum of the last two terms \(Z_{\text{SN}}[\dot{M}(\text{SN} \rightarrow \text{ISM} \rightarrow \text{out}) + \dot{M}(\text{SN} \rightarrow \text{out})]\) in Equation (21), corresponding to the “cumulative” SN-origin metal flux.

While not directly calculated, the “instantaneous” SN-origin metal flux, \(Z_{\text{SN}}\dot{M}(\text{SN} \rightarrow \text{out})\), is of great interest to quantify how much of injected metals go promptly into outflows and how enriched the outflow is compared to the ISM. We use the following procedure to estimate this quantity. Theoretically, the
instantaneous metallicity of ISM-origin outflows is
\[ Z_{\text{ISM}} \equiv \frac{Z_{\text{ISM},0} M_{\text{ISM}} (\text{ISM} \rightarrow \text{out}) + Z_{\text{SN}} M_{\text{SN}} (\text{SN} \rightarrow \text{ISM} \rightarrow \text{out})}{M (\text{ISM} \rightarrow \text{out}) + M (\text{SN} \rightarrow \text{ISM} \rightarrow \text{out})} \tag{22} \]
while the mean metallicity of outflows,
\[ Z = \frac{M}{M} , \tag{23} \]
is directly measured in the simulation using mass and total metal scalar fluxes at specified \( z \). Combining with Equations (20) and (21), we obtain the instantaneous SN-origin mass outflow rate
\[ \dot{M} (\text{SN} \rightarrow \text{out}) = \frac{Z - Z_{\text{ISM}}}{Z_{\text{SN}} - Z_{\text{ISM}}} M \equiv \dot{f}_{\text{SN}} M \tag{24} \]
and the instantaneous SN-origin metal outflow rate
\[ \dot{M}_{Z} (\text{SN} \rightarrow \text{out}) = \frac{Z_{\text{SN}}}{Z_{\text{SN}} - \dot{f}_{\text{SN}}} \dot{M} (\text{SN} \rightarrow \text{out}) = \frac{Z_{\text{SN}}}{\dot{f}_{\text{SN}}} \dot{M}_{Z} \tag{25} \]
Equations (24) and (25) define the instantaneous SN-origin mass and metal fractions in outflows, \( \dot{f}_{\text{SN}} \) and \( \dot{f}_{\text{SN}} \), respectively. In the rest of the paper, we will use the superscript “SN” to refer to the instantaneous SN-origin component, e.g., \( M_{\text{out}}^\text{SN} \) and \( M_{Z\text{out}}^\text{SN} \) for \( M (\text{SN} \rightarrow \text{out}) \) and \( M_{Z} (\text{SN} \rightarrow \text{out}) \), respectively, and \( \bar{f}_{\text{SN}}^\text{M} \) and \( \bar{f}_{\text{SN}}^Z \) for the corresponding mass and metal fluxes.

As a proxy for the instantaneous metallicity of ISM-origin outflows in Equation (22), we use the instantaneous metallicity of the ISM itself. In practice, in the simulation we measure the ISM metallicity based on the cool-phase gas within \(| z | < 50 \) pc; this defines \( Z_{\text{ISM}} \) (we find no strong variation with different thickness used in this definition if smaller than the scale height). For a given \( Z_{\text{ISM}} \), we use the phase-separated mean outflow metallicity (\( Z_{\text{ph}}^\text{M} \)) and mass outflow rate (\( \dot{M}_{\text{ph}}^\text{M} \)) to obtain \( \dot{f}_{\text{SN}}^\text{M} \) and \( \bar{f}_{\text{SN}}^\text{Z} \) phase by phase. Note that our definition for \( Z_{\text{ISM}} \) is not a perfect tracer of the instantaneous metallicity in ISM-origin outflows, so that occasionally we get negative \( \dot{f}_{\text{SN}} \); we simply set it to zero in such occasions. This occurs only for the cool outflow at \(| z | = H \) (at most 20% of the time). R16 is the only exception, where the genuine cool ISM is easily pushed out to large distance so that \( Z_{\text{cool}} \leq Z_{\text{ISM}} \) for most snapshots at all heights (up to 80% of the time). For this reason, we exclude R16 in analysis regarding SN-origin metals of cool outflows (e.g., Table 5 and Figure 8). For the hot gas, \( \dot{f}_{\text{SN}}^\text{M} \) is always positive.

Given instantaneous outflow and ISM metallicities, we obtain the instantaneous outflow enrichment factor for each phase “ph”
\[ \zeta_{\text{ph}} = \frac{Z_{\text{ph}}}{Z_{\text{ISM}}} . \tag{26} \]

Figure 5 shows, from the top to bottom, (a) the mean metallicity of outflow along with the instantaneous ISM metallicity (solid dashed), (b-c) the fractions of instantaneous SN-origin mass and metal in the outflow, and (d) the instantaneous outflow enrichment factor for R4. The ISM is gradually enriched, and cool outflows consist mostly of the pre-enriched ISM (\( \zeta_{\text{cool}} \sim 1 \)). The hot outflow is more metal enriched than cooler components and the ISM by a factor of 1.5–2. The contribution of recent-SN-origin metals to the outflowing metal flux is \( \sim 30\%–60\% \) in the hot outflow and \( \sim 10\% \) in the cooler components.

4.3. Loading Factors

We now calculate loading factors; the ratios of outgoing mass, momentum, energy, and metal to mass locked into stars and momentum, energy, and metal injected by SNs. With the definition of the reference outflow fluxes in Equation (15), we get outflow loading factors simply by
\[ \eta_q = \frac{\bar{f}_q}{\bar{f}_{\text{SN}}} , \tag{27} \]
where \( q = M, p, E, \) and \( Z \). The definitions of \( \eta_M \) and \( \eta_E \) are identical to the conventional definition (e.g., Chevalier & Clegg 1985).

In principle, Equation (27) can give instantaneous loading factors as a function of time, but care needs to be taken with this. The quantity of interest is the outflow rate normalized by the injection (or star formation) rate that is responsible for the outflow. The injection (or star formation) generally occurs near the disk midplane, while the outflow is measured at a certain height above the midplane. There is inevitably a time delay between injection and outflow rates. Therefore, instantaneous loading factors measured by Equation (27) can be misleading regarding the physical impact of stellar feedback. This issue is more serious when (1) star formation is more bursty and the SFR rather than SNR is used for the reference flux;\(^{14}\) and (2) the distance between locations where feedback injection (or star formation) and outflow rates are measured is larger. One would need to either carefully model the reference fluxes, including determination of an appropriate time delay in computing instantaneous loading factors, or else report time-averaged loading factors with averaging timescale longer than the time delay and timescale of the feedback cycle (e.g., Muratov et al. 2015).

Taking advantage of the long duration of our simulation suite (covering a few feedback/outflow cycles), we report the ratio of time-averaged flux to time-averaged reference flux as time-averaged loading factors. For model R4, Figure 6 shows the vertical profiles of the time-averaged loading factors, as well as temporal variation ranges. We note that in contrast to Figure 4, rather than the total metal loading factor we now show the instantaneous SN-origin metal loading factor \( \eta^Z \) (using Equation (25)) in the bottom panel of Figure 6. We also note that the reference fluxes are not height dependent, so that Figure 6 essentially shows rescaled flux profiles. We plot as symbols with error bars the mean and standard deviation of loading factors measured at the instantaneous \( H \) (circle) and \( 2H \)

\(^{14}\) For example, Martizzi (2020) recently reported a large instantaneous mass loading factor (\( \sim 100 \)), which they suggested was due to clustered star formation under self-gravity. However, from Figures 7 and 10 in Martizzi (2020), the peak in the instantaneous mass loading factor from Model S100_WSG at \( t/t_{\text{dyn}} \sim 3 \) occurs at both a maximum of the outflow rate and a minimum of the SFR. If one simply reads off the peaks of both outflow and SFRs and takes the ratio, the mass loading factor is 0.1, comparable to their non-self-gravitating (nonclustering) model. The drop in star formation (after an initial big burst) is the major reason for the high instantaneous mass loading factor.
square), which are generally in good agreement with the values from the time-averaged profiles at the mean $H$ and $2H$. In Appendix B, we make use of the time-delayed reference fluxes to find the mean time delay and calculate instantaneous loading factors. The time-averaged profiles of the instantaneous loading factors are almost identical to Figure 6, providing reassurance of the robustness of mean loading factors we report here.

In Figure 6, we see a steep drop of all loading factors of the cool phase from the midplane to $H$. The hot (and intermediate) phase loading factors peak at $\sim 50$ pc (most SNe explode below this height). Above $|z| \sim 50$ pc, outflow fluxes gradually drop. The decreasing trend is moderated above $H$ but is still significant in cooler phases for mass and hotter phases for energy. The mass loading factor of the cool phase decreases with $|z|$ as lower velocity components drop out (see Figure 7 and KO18). The energy loading factor of the hot phase decreases from a maximum slightly above the midplane as some of the energy (mostly thermal) in the hot gas is transferred to cooler phases, from which it is quickly radiated away (Vijayan et al. 2020). The intermediate phase is subdominant for all loading factors at all heights (except R16).

The momentum loading factor of the sum of the cool and hot components near the midplane is $\eta_p \sim 0.5$, implying that SNRs have built-up momentum exceeding the initial ejecta momentum (which would yield $\eta_p \sim 0.1$). We note that $\eta_p$ is not as large as unity since the terminal momentum per SN from clustered SNe is generally smaller than $p_{\text{ref}}$ from a single SN,
especially when SN events are nearly continuous and blast waves become subsonic in the hot ISM before reaching the cool shell (Kim et al. 2017a; El-Badry et al. 2019). Also, a portion of the injected SN momentum flux is converted to magnetic stresses, and near the midplane these are comparable to the vertical kinetic momentum flux. The momentum flux also decreases as a function of \( z \), especially for the cool component, since it must contribute support against the weight of the ISM (the thermal plus turbulent pressure is approximately twice \( \rho \) \( c_s^2 \)), allowing for \( v_{\text{out}} < 0 \). At \( |z| = H \) and above, the leftover kinetic vertical momentum flux is only 10% of the reference momentum flux. This is generally true except in R16, in which fluxes are all dominated by the cool component and SN events are more or less discrete.

In Table 3, we provide the mean values over \( 0.5 < t/t_{\text{eis}} < 1.5 \) of the measured fluxes and loading factors of all models and phases at \( |z| = H \).

4.4. Characteristic Velocities and Metal Properties

Figure 7 plots time-averaged vertical profiles of additional quantities of interest, including (a) outflow velocity, (b) Bernoulli velocity, (c) metallicity, and (d) metal enrichment factor for model R4. The characteristic outflow velocity is defined as

\[
\tau_{\text{out}, \text{ph}}(z; t) \equiv \frac{F_{p, \text{ph}}^{\text{kin}}(z; t)}{F_{M, \text{ph}}^{\text{kin}}(z; t)},
\]

where \( F_{p, \text{ph}}^{\text{kin}} \) is the kinetic component of momentum flux defined by only the first term of Equation (12). The Bernoulli
## Table 3
Time-averaged Fluxes and Loading Factors at \(|z| = H\)

| Model (1) | Phase (2) | \(F_M (3)\) | \(F_P (4)\) | \(F_E (5)\) | \(F_{SN} (6)\) | \(\eta_M (7)\) | \(\eta_P (8)\) | \(\eta_E (9)\) | \(\eta_F (10)\) | \(\eta_{SN} (11)\) | \(\eta_{SN}^{BN} (12)\) |
|-----------|-----------|-------------|-------------|-------------|-------------|-------------|-------------|-------------|-------------|-------------|-------------|
| R2        | cool      | 0.75        | 51          | 7.3 \times 10^{46} | 2.9 \times 10^{-2} | 3.2 \times 10^{-2} | 0.69        | 3.5 \times 10^{-2} | 6.4 \times 10^{-3} | 1.3          | 0.14        |
|           | int       | 6.3 \times 10^{-2} | 10          | 2.8 \times 10^{46} | 2.6 \times 10^{-3} | 5.6 \times 10^{-4} | 5.8 \times 10^{-2} | 7.1 \times 10^{-3} | 2.5 \times 10^{-3} | 0.11         | 2.5 \times 10^{-2} |
|           | hot       | 0.13        | 1.4 \times 10^3 | 2.8 \times 10^{48} | 9.6 \times 10^{-3} | 6.2 \times 10^{-3} | 0.12        | 0.10         | 0.24         | 0.42         | 0.27        |
| R4        | cool      | 0.27        | 12          | 1.1 \times 10^{46} | 8.3 \times 10^{-3} | 4.4 \times 10^{-4} | 2.2         | 7.7 \times 10^{-2} | 8.2 \times 10^{-3} | 3.3          | 0.17        |
|           | int       | 1.4 \times 10^{-2} | 1.9          | 4.2 \times 10^{45} | 4.8 \times 10^{-4} | 7.2 \times 10^{-5} | 0.12        | 1.2 \times 10^{-2} | 3.3 \times 10^{-3} | 0.19         | 2.8 \times 10^{-2} |
|           | hot       | 2.7 \times 10^{-2} | 19          | 2.2 \times 10^{47} | 1.3 \times 10^{-3} | 6.0 \times 10^{-4} | 0.22        | 0.12         | 0.17         | 0.51         | 0.23        |
| R8        | cool      | 3.3 \times 10^{-2} | 0.79        | 4.4 \times 10^{44} | 7.2 \times 10^{-4} | 2.1 \times 10^{-5} | 6.4         | 0.12         | 8.2 \times 10^{-3} | 6.7          | 0.20        |
|           | int       | 1.3 \times 10^{-3} | 0.12        | 2.3 \times 10^{44} | 3.0 \times 10^{-5} | 2.9 \times 10^{-6} | 0.25        | 1.8 \times 10^{-2} | 4.3 \times 10^{-3} | 0.28         | 2.7 \times 10^{-2} |
|           | hot       | 1.3 \times 10^{-3} | 0.67        | 5.5 \times 10^{45} | 4.1 \times 10^{-5} | 1.5 \times 10^{-5} | 0.26        | 0.10         | 0.10         | 0.39         | 0.14        |
| R16       | cool      | 5.5 \times 10^{-3} | 8.5 \times 10^{-2} | 2.3 \times 10^{43} | 1.1 \times 10^{-4} | 3.3 \times 10^{-6} | 56         | 0.67         | 2.2 \times 10^{-2} | 54           | 1.6 \times 10^{-3} |
|           | int       | 3.6 \times 10^{-3} | 2.9 \times 10^{-3} | 3.8 \times 10^{42} | 7.8 \times 10^{-7} | 5.3 \times 10^{-8} | 0.37        | 2.3 \times 10^{-2} | 3.8 \times 10^{-3} | 0.39         | 2.6 \times 10^{-2} |
|           | hot       | 1.4 \times 10^{-5} | 9.3 \times 10^{-3} | 6.1 \times 10^{43} | 4.5 \times 10^{-7} | 1.8 \times 10^{-7} | 0.15        | 7.3 \times 10^{-2} | 6.0 \times 10^{-2} | 0.22         | 8.7 \times 10^{-2} |
| LGR2      | cool      | 0.55        | 27          | 2.8 \times 10^{46} | 1.8 \times 10^{-2} | 1.5 \times 10^{-3} | 1.2         | 4.2 \times 10^{-2} | 5.7 \times 10^{-3} | 1.9          | 0.15        |
|           | int       | 2.6 \times 10^{-2} | 3.6          | 8.9 \times 10^{45} | 9.7 \times 10^{-4} | 1.9 \times 10^{-4} | 5.4 \times 10^{-2} | 5.7 \times 10^{-3} | 1.8 \times 10^{-3} | 9.7 \times 10^{-2} | 1.9 \times 10^{-2} |
|           | hot       | 5.4 \times 10^{-2} | 48          | 6.7 \times 10^{47} | 3.2 \times 10^{-3} | 1.8 \times 10^{-3} | 0.11        | 7.6 \times 10^{-2} | 0.14         | 0.33         | 0.18        |
| LGR4      | cool      | 0.46        | 14          | 8.4 \times 10^{45} | 1.2 \times 10^{-2} | 2.1 \times 10^{-4} | 5.1         | 0.12         | 9.0 \times 10^{-3} | 6.3          | 0.11        |
|           | int       | 1.0 \times 10^{-2} | 1.2          | 2.5 \times 10^{45} | 3.0 \times 10^{-4} | 3.7 \times 10^{-5} | 0.11        | 1.0 \times 10^{-2} | 2.7 \times 10^{-3} | 0.16         | 2.0 \times 10^{-2} |
|           | hot       | 1.5 \times 10^{-2} | 10          | 1.0 \times 10^{47} | 6.5 \times 10^{-4} | 2.8 \times 10^{-4} | 0.17        | 8.4 \times 10^{-2} | 0.11         | 0.34         | 0.15        |
| LGR8      | cool      | 4.0 \times 10^{-2} | 0.85        | 3.6 \times 10^{44} | 8.6 \times 10^{-4} | 7.8 \times 10^{-6} | 13         | 0.20         | 1.1 \times 10^{-2} | 13           | 0.12        |
|           | int       | 7.3 \times 10^{-4} | 7.2 \times 10^{-2} | 1.3 \times 10^{44} | 1.7 \times 10^{-5} | 1.4 \times 10^{-5} | 0.23        | 1.7 \times 10^{-2} | 4.0 \times 10^{-3} | 0.26         | 2.2 \times 10^{-2} |
|           | hot       | 8.8 \times 10^{-4} | 0.43        | 3.3 \times 10^{35} | 2.7 \times 10^{-5} | 8.5 \times 10^{-6} | 0.28        | 0.10         | 9.9 \times 10^{-2} | 0.41         | 0.13        |

Note: Columns (3)–(7) are time-averaged outflow fluxes defined by Equation (10) for (3) mass flux in \(M_e/\text{km s}^{-1}\) (Equation (11)), (4) momentum flux in \((M_e \times \text{km s}^{-1})/\text{pc}^2 \text{yr}^{-1}\) (Equation (12)), (5) energy flux in \(\text{erg}/\text{km}^2 \text{pc}^2 \text{yr}\) (Equation (13)), (6) total metal flux in \(M_e/\text{pc}^2 \text{yr}\) (Equation (14)), and (7) SN-origin metal flux in \(M_e/\text{pc}^2 \text{yr}\) (Equation (25)). Columns (8)–(12) are time-averaged loading factors (dimensionless) defined by Equations (27) and (15) with (8) for mass, (9) for momentum, (10) for energy, and (11) for total metal and (12) SN-origin metal. Time averages are taken over \(0.5 \text{t}_{\text{sh}} < t < 1.5 \text{t}_{\text{sh}}\). The data for additional tables at different heights \(|z| = 2H, 500 \text{ pc}, \text{ and } 1 \text{ kpc}\), as well as standard deviations, are available at doi:10.5281/zenodo.3872049.
For an adiabatic steady flow, the Bernoulli velocity at $z$ must exceed the escape speed $v_{\text{esc}} = \left[ 2 (\Phi(z_{\text{max}}) - \Phi(z)) \right]^{1/2}$ in order for the flow to reach $z_{\text{max}}$; this criterion also applies for a completely cold ballistic flow, where $v_{\text{g}} \rightarrow v$.

The outflow velocity of cool outflows is as low as $\sim 30 \, \text{km} \, \text{s}^{-1}$ near the midplane and increases as the outflow moves farther away, reaching as high as $\sim 100 \, \text{km} \, \text{s}^{-1}$ near the simulation boundaries. The increasing trend of the outflow velocity of the cool phase with $|z|$ seen in Figure 7(a) is often interpreted as an acceleration, but the mean outflow velocity can also increase as low-velocity gas drops out. The former is more important near the midplane $|z| < H$, where actual acceleration of the cool phase by superbubble expansion is occurring, but the latter dominates the trend at higher altitudes (e.g., KO18, Vijayan et al. 2020). Indeed, panel (b) of Figure 6 shows a trend of steadily decreasing momentum flux with $|z|$, due to the dropout of low-velocity gas. In the extraplanar region $|z| > H$, some acceleration of cool outflows occurs owing to the hot—cool interaction, which helps to maintain high-velocity tails of cool outflows (Vijayan et al. 2020), but this is not the dominant reason for the increasing trend of outflow velocity. We also note that cooling of intermediate-temperature gas is preferentially at low velocity and adds to the cool gas inflow; cooling of intermediate-temperature gas has minimal impact on the momentum transfer to the cool phase (see Vijayan et al. 2020).

For model R4, the escape velocity from the box relative to $|z| = H = 340 \, \text{pc}$ (where we tabulate $v_{\text{g}}$) is $v_{\text{esc}} = 154 \, \text{km} \, \text{s}^{-1}$. Given the low mean outflow velocity of cool outflows, it is evident that the majority of cool-phase (and intermediate-phase) outflows cannot travel far from the disk midplane and escape the simulation domain. This is also clearly demonstrated by the steep decrease of mass loading factor as a function of $z$ in Figure 6(a).

The outflow velocity of hot outflows, in contrast, is higher than the escape velocity of the system, as clearly illustrated in Figure 3. The Bernoulli velocity is much larger than the outflow velocity as the thermal term dominates, implying the possibility of further acceleration of hot outflows. In our simulations, the outflow velocity for the hot gas flattens out above $|z| > H$ or $2H$, reaching $\frac{v_{\text{out}}}{v_{\text{esc}}} \sim 250 \, \text{km} \, \text{s}^{-1}$ for model R4. This flattening is mainly due to the limited volume of the local box simulations. When a volume much larger than the source region is available, hot outflows expand and increase outflow velocity at the expense of thermal energy (e.g., Chevalier & Clegg 1985). In order for a hot flow to fully accelerate, the simulation box must be large compared to the source region, so that the streamlines can open and transition through a sonic point before reaching the boundaries (e.g., Fielding et al. 2017; Schneider & Robertson 2018; Schneider et al. 2020), which generally does not occur when there is distributed star formation in a local box (e.g., Martizzi et al. 2016).
Since the asymptotic velocity is \( v = (v_{\text{g}}^2 - v_{\text{esc}}^2)^{1/2} \) for an adiabatic wind, the Bernoulli velocity can be used as a proxy for the terminal velocity that the hot gas would reach in the case that \( v_{\text{g}} \gg v_{\text{esc}} \). In Figure 7(b), the Bernoulli velocity of hot outflow decreases with \( z \), from \( v_{\text{g,hot}}(H) \approx 820 \text{ km s}^{-1} \) to \( v_{\text{g,hot}}(2H) \approx 660 \text{ km s}^{-1} \) to \( v_{\text{g,hot}}(L_c/2) \approx 490 \text{ km s}^{-1} \), while the escape velocity decreases from \( v_{\text{esc}}(H) = 154 \text{ km s}^{-1} \) to \( v_{\text{esc}}(2H) = 137 \text{ km s}^{-1} \). Since only the combination \( v_{\text{g}}^2 - v_{\text{esc}}^2 \) is expected to be conserved in an adiabatic flow, the decrease of \( v_{\text{g}} \) with distance is due in part to the decrease of \( v_{\text{esc}} \). This effect is small when \( v_{\text{g}} \gg v_{\text{esc}} \). Within the main body of the disk \( (z < H \text{ or } 2H) \), a decrease in \( v_{\text{g}} \) is also expected since the strong interaction between hot and cool components transfers energy from hot to cooler gas. After the hot gas emerges into the extraplanar region, the interaction between phases is reduced, but there is still substantial loss of energy flux from the hot component owing to interaction with cool fountain flows populated by previous events (Vijayan et al. 2020). Even with these losses, the Bernoulli velocity of the hot outflow in all models is large enough \( (>600 \text{ km s}^{-1} \) at \( z = 2H \)) that the hot gas could be expected to travel far out into the CGM.

Figure 7(c) plots the mean metallicity of outflows. As shown in Figure 5(a), the metallicity of outflows (and the ISM) gradually increases over time. SNe inject metals mainly near the midplane. As hot, metal-enriched bubbles expand and mix into surrounding cooler gas, the metallicity of the hotter/cooler component decreases/increases as outflows travel farther. The mean metallicity in each phase significantly changes as a function of \( z \) up to \( z = 2H \); again indicating active interaction and mixing between phases within \( z < 2H \). For \( z \in (H, 2H) \), the hot and cool outflows are respectively \( \sim 50\% \) and \( 10\% - 20\% \) more metal enriched than the ISM near the midplane (Figure 7(d)).

In Table 4, we provide the mean values of the mass-weighted outflow velocity, Bernoulli velocity, mean metallicity, and enrichment factors of all models and phases at \( z = H \), averaged over 0.5 \(< t/t_{\text{orb}} < 1.5 \).

### 5. Scaling Relations

In this section, we systematically investigate the dependence of outflow characteristics (loading factors, metal properties, and outflow velocities) on a variety of galactic properties in our simulations, including SFR surface density \( (\Sigma_{\text{SFR}}) \), gas surface density \( (\Sigma_{\text{gas}}) \), midplane gas number density \( (n_{\text{mid}}) \), midplane total pressure \( (P_{\text{mid}}) \), gas weight \( (VV) \), and gas depletion time \( (t_{\text{dep}}) \). At any time, \( \Sigma_{\text{SFR}} \) is calculated from the total mass of star cluster particles with age younger than \( t_{\text{bin}} \) such that

\[
\Sigma_{\text{SFR},t_{\text{bin}}} = \frac{M_{\text{gas}}(t_{\text{age}} < t_{\text{bin}})}{t_{\text{bin}}L_y L_y}.
\]

As a default, we use \( t_{\text{bin}} = t_{\text{evol}} = 40 \text{ Myr} \), corresponding to the SFR definition that best traces the SN rates used in the reference flux calculations, but we also explored different \( t_{\text{bin}} = 10 \text{ Myr} \) and 100 Myr. \( \Sigma_{\text{gas}} = M_{\text{gas}}/(L_xL_y) \) is directly calculated from the total gas mass divided by horizontal area. Midplane averages are computed by taking averages in two horizontal slices at \( z = \pm \Delta z/2 \), with \( n_{\text{mid}} \) and \( P_{\text{mid}} \) defined using volume-averaged number density and total pressure (including turbulent, thermal, and magnetic terms) just for cool gas. \( VV \) is obtained by directly integrating \( \rho \delta \delta'/dz \) for cool gas from the top or bottom of the simulation domain to the midplane and averaging the two values. The depletion time \( t_{\text{dep}} = \Sigma_{\text{gas}}/\Sigma_{\text{SFR},t_{\text{bin}}} \) for \( t_{\text{bin}} = 40 \text{ Myr} \). Here we will present dependencies on galactic properties as scaling relations for cool and hot phase loading factors, characteristic velocities, and metal enrichment measured at \( z = H \). We also have fit scaling relations at different heights, and these results are available at doi:10.5281/zenodo.3872049 (see Jupyter notebook)

We generally find smaller intrinsic scatter and better correlation at \( z = H \) and 2H than at fixed heights \( z = 500 \text{ pc} \) and 1 kpc. Because the intermediate phase is subdominant, we do not include these results in this section, but the data are available at doi:10.5281/zenodo.3872049.

To quantify scaling relations between two variables, we report linear regression results in log-log space. We first construct time series of quantities of interest with 0.01\( t_{\text{orb}} \) interval over \( 0.5 < t/t_{\text{orb}} < 1.5 \) for each model. We then perform bootstrap resampling 500 times with a sample size of 10 (we find typical autocorrelation timescales of time series \( t_{\text{corr}}/t_{\text{orb}} \in (0.05, 0.1) \)) to obtain the mean (\( \bar{q} \)) and its error (\( \delta q/\bar{q} \)). We feed in log of the mean (\( \log \bar{q} \)) and error (\( \delta q/\bar{q} \)) for linear regression using a Python version of the linmix package.16 This is a widely tested Bayesian estimator for linear regression (Kelly 2007) to derive posterior distributions of intercept \( \alpha \) and slope \( \beta \), as well as intrinsic scatter \( \sigma_{\text{int}} \) and Pearson correlation coefficient \( \rho \).

#### 5.1. Loading Factors with SFRs

Figure 8 shows scaling relations of mass, momentum, energy, and SN-origin metal loading factors measured at \( H \) as a function of \( \Sigma_{\text{SFR}} \) for cool (left) and hot (right) outflows. We present the mean and error of measured quantities from bootstrapping as symbols and error bars, which we use for the fitting, along with small points denoting time evolution of each model over \( 0.5 < t/t_{\text{orb}} < 1.5 \) with a sampling interval of 0.01\( t_{\text{orb}} \). In each panel, the solid line and shaded regions denote the median and 68\% and 95\% confidence intervals of model posterior distributions; the median and 16th/84th percentile values of the intercept \( \alpha \), slope \( \beta \), intrinsic scatter \( \sigma_{\text{int}} \) and Pearson correlation coefficient \( \rho \) are shown in the box of each panel (also shown in Table 5).

The hot mass loading factors are nearly flat with \( \eta_{\text{H,hot}} \sim 0.1-0.2 \) over a wide range in \( \Sigma_{\text{SFR}} \). This level of the hot gas loading is consistent with what has been reported in other simulations (e.g., Li et al. 2017; Li & Bryan 2020) and with the expectation from superbubble breakout after shell formation (Kim et al. 2017a). The hot gas energy loading factors \( \eta_{\text{E,hot}} \) show a weakly increasing trend from 0.06 to 0.25 with \( \Sigma_{\text{SFR}} \) and are larger than \( \eta_{\text{E,cool}} \) by more than an order of magnitude. In general, the models with higher SFR have greater temporal and spatial correlation of star formation (and SNe), providing a potential explanation for the enhancement of energy loading factor. However, the effect is less dramatic than suggested by previous idealized numerical simulations (Fielding et al. 2018). This is partly because we are reporting time-averaged loading factors (averaging over both high and low states) and partly because our self-consistent

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15 https://github.com/changgoowinds/figset/blob/paper1/tables/Example_scripts.ipynb
16 https://github.com/jmeyers314/linmix
simulations always have fountain flow gas at high altitudes, with which hot gas must interact. In addition, the larger horizontal velocity dispersions at higher SFR tend to close off chimneys. Thus, even though there is a burst of star formation that creates a superbubble, the energy loading is reduced below what it would be if the superbubble were to vent into an almost-vacuum region.

The cool mass loading factors $\eta_{M,\text{cool}}$ decrease steeply with $\Sigma_{\text{SFR}}$, with values ranging from 100 to 1. However, it is noteworthy that much of the cool gas is at low velocities ($v_{\text{out,cool}} \sim 10$–100 km s$^{-1}$; see Table 4), as evidenced by the low energy-loading factor (see also KO18). Therefore, the high mass loading factor of the cool phase at $|z| = H$ shown here does not immediately imply heavily mass-loaded winds at large distances. Indeed, the mass loading factor in model R4 drops by a factor of 3 from $|z| = H$ to $|z| = 2H$ and keeps decreasing as a function of $|z|$ (see Figure 6 and also KO18).

In our simulation suite, most of the mass in cool outflows cannot reach the vertical boundary of the simulation box and falls back toward the midplane (see Figure 3). It is still possible to
### Table 5

| X (1) | Y | α (3) | β (4) | Cov(α, β) | σ_m (5) | ρ (7) | X (8) | Y | α (9) | β (10) | Cov(α, β) | σ_m (11) | ρ (12) |
|------|---|-------|------|----------|---------|------|------|---|---|-------|------|----------|---------|------|
|      |   |       |      |          |         |      |      |   |   |       |      |          |         |      |
|        |   |       |      |          |         |      |      |   |   |       |      |          |         |      |
|        |   |       |      |          |         |      |      |   |   |       |      |          |         |      |
|        |   |       |      |          |         |      |      |   |   |       |      |          |         |      |
|        |   |       |      |          |         |      |      |   |   |       |      |          |         |      |
|        |   |       |      |          |         |      |      |   |   |       |      |          |         |      |
|        |   |       |      |          |         |      |      |   |   |       |      |          |         |      |
|        |   |       |      |          |         |      |      |   |   |       |      |          |         |      |
|        |   |       |      |          |         |      |      |   |   |       |      |          |         |      |

### Note

Linear regression results for logX and logY. We exclude R16 for fitting of \(\eta_{\text{SN}}\). The values given for the intercept, slope, intrinsic scatter \(\sigma_{\text{int}}\) and Pearson correlation coefficient \(\rho\) are the median and interval containing 68% of the estimates over the posterior distributions. Covariance of \(\sigma_{\text{int}}\) and \(\rho\) is given in Columns (5) and (10). The data at different heights along with Python scripts for fitting are available at doi:10.5281/zenodo.3872049.

anticipate a higher mass loading factor at large distances (e.g., 0.1–1 virial radius) in dwarf galaxies that have a shallower gravitational potential, as reported in cosmological zoom-in (Martizzi et al. 2015) or isolated galaxy simulations (Hu 2019). We refrain from extrapolating our results to that regime since those outflows may consist of both directly launched cool outflows and swept-up CGM driven out by energy delivered by hot outflows.

The momentum loading factors in the cool gas at \(|z| = H\) decrease from \(\eta_{\text{p,cool}} \sim 0.7\) to 0.4 with increasing \(\Sigma_{\text{SFR}}\). When
combined with the nearly constant $\eta_p, \text{hot} \sim 0.1$ and decreasing trend of $\eta_s$ as a function of $z$ in general (see Figure 6), this implies that most of the vertical momentum injection from SN feedback goes into the bulk of the ISM in the disk, rather than escaping from galaxies. Further analysis of the momentum injection to the ISM from SN feedback, quantifying its contribution to supporting the gravitational weight of the disk and regulating SFRs, will be given in a separate paper (E. Ostriker & C.-G. Kim 2020, in preparation; see also Kim et al. 2011, 2013; Kim & Ostriker 2015b). Consistent with our previous result in KO18 (see also Li & Bryan 2020), we find that the energy loading factor of the cool gas is significantly lower than in the hot gas, $\eta_{E, \text{cool}} \sim 0.02–0.005$, decreasing with increasing $\Sigma_{\text{SFR}}$.

A key conclusion from our simulation suite is that energy is carried by hot outflows while mass is carried by cool outflows. Including two distinct wind components is therefore crucial in any physically motivated wind model.

### 5.2. Dependence of Loading Factors on Galactic Properties

Figures 9 and 10 show $\eta_{M, \text{cool}}$ and $\eta_{E, \text{hot}}$ as a function of different galactic conditions, including $\Sigma_{\text{SFR}}$ with different $\tau_{\text{bin}}$, $\Sigma_{\text{gas}}$, $n_{\text{mid}}$, $P_{\text{mid}}$, $\mathcal{W}$, and $t_{\text{dep}}$ with $\tau_{\text{bin}} = 40$ Myr. These parameters are chosen both because they represent important physical properties of the ISM and because we expect outflows to correlate with them. These parameters are also quantities that can be estimated in large-volume cosmological simulations or SAMs and therefore would be available as inputs to a subgrid model for wind launching. The level of $\Sigma_{\text{SFR}}$ (which is not imposed but obtained self-consistently in each simulation) sets the overall strength of feedback, while $\Sigma_{\text{gas}}$ and $n_{\text{mid}}$ characterize the conditions that affect superbubble propagation and breakout, as well as transfer of momentum and energy to the bulk ISM. The values of $P_{\text{mid}}$ and $\mathcal{W}$ are related to each other and to $\Sigma_{\text{SFR}}$ through self-regulation. These also reflect the vertical gravitational field and therefore encode oscillation timescales that control fountain flows that limit gas escape. The value of $t_{\text{dep}} = \Sigma_{\text{gas}}/\Sigma_{\text{SFR}}$ represents a local secular evolutionary timescale.

In the first three panels, (a)–(c), of Figures 9 and 10, we compare scaling relations for three different choices of the averaging timescale for $\Sigma_{\text{SFR}}$. These values, $\tau_{\text{bin}} = 10$, 40, and 100 Myr, are rough proxies for different observational tracers. As expected, $\Sigma_{\text{SFR}}$ traced by younger star clusters (e.g., $\Sigma_{\text{SFR,10}}$) exhibits larger-amplitude fluctuations. Nevertheless, the scaling relations from time-averaged points with all $\tau_{\text{bin}}$ choices are consistent with each other.

Panels (d)–(h) of Figures 9 and 10 show scaling relations of $\eta_{M, \text{cool}}$ and $\eta_{E, \text{hot}}$ with respect to $\Sigma_{\text{gas}}$, $n_{\text{mid}}$, $P_{\text{mid}}$, $\mathcal{W}$, and $t_{\text{dep}}$. Overall, we do not find particularly better correlation with one parameter over another (except that $\Sigma_{\text{gas}}$ has poorer correlation). This is mainly because the quantities are not physically independent but mutually connected through self-regulation. Fundamentally, the SFR surface density is self-regulated to provide the vertical pressure support through feedback that is required by the gas weight $\Sigma_{\text{SFR}} \propto P_{\text{mid}} \approx \mathcal{W}$ (Ostriker et al. 2010; Kim et al. 2011; Ostriker & Shetty 2011) with near-linear relationships demonstrated in both simulations (E. Ostriker & C.-G. Kim 2020, in preparation; see also Shetty & Ostriker 2012; Kim et al. 2013; Kim & Ostriker 2015b) and observations (Herrera-Camus et al. 2017; Sun et al. 2020). All three of these quantities therefore are fundamental measures of the feedback strength, while including the local vertical gravity and gas density implicitly/explicitly. The midplane pressure and the weight are the same on average, but their instantaneous response to feedback is different; the midplane pressure responds more immediately and directly to SN rates ($\propto \Sigma_{\text{SFR,40}}$) and FUV luminosity ($\propto \Sigma_{\text{SFR,10}}$), while the weight varies only indirectly through the change of gas scale height (or velocity dispersion). The temporal variations in $P_{\text{mid}}$ and $\mathcal{W}$ are thus similar to those in $\Sigma_{\text{SFR}}$ with shorter and longer averages, respectively, so that the scatter in the points in panels (f) and (g) is more or less similar to panels (a)/(b) and (c), respectively.

The scaling with $\Sigma_{\text{gas}}$ (panel (d)) is related to the scaling with gas weight. If the external gravity dominates the weight, $\mathcal{W} \approx \Sigma_{\text{gas}} \sigma (2G\rho_{\text{ad}})^{1/2}$, where $\rho_{\text{ad}} \equiv \rho_{\text{S}}/(2\Sigma_{\infty})$ is the midplane density of stars and dark matter; however, a large range of $\Sigma_{\text{SFR}} \propto \mathcal{W}$ is possible at a given $\Sigma_{\text{gas}}$. As a consequence, correlation with $\Sigma_{\text{gas}}$ is indeed slightly worse than other parameters considered, including $\Sigma_{\text{SFR}}$, $P_{\text{mid}}$, and $\mathcal{W}$, judging from the larger intrinsic scatter and the smaller Pearson correlation coefficient derived by linear regression. A wider parameter space survey and more experiments with extreme combinations between gas and gravity parameters would help to uncover which properties are the most fundamental in setting the loading factors.

The scaling with $n_{\text{mid}}$ (panel (e)) is a measure of cooling in the ISM ($E_{\text{cool}} \sim n^2$) and is also related to the scaling with midplane pressure, since $P_{\text{mid}} \sigma_{s,\text{eff}}^2 \approx \rho_{\text{mid}}$. Over more than three orders of magnitude variation in $P_{\text{mid}}$ covered by our simulation suite, the effective vertical velocity dispersion $\sigma_{s,\text{eff}}$ increases by no more than a factor of 3 from the lowest to the highest $\Sigma_{\text{SFR}}$ and $P_{\text{mid}}$ (see Table 2; see also Joung et al. 2009). Therefore, panels (e) and (f) are similar.

Finally, the scatter in the relation for $t_{\text{dep}}$ (panel (h)) in each model simply arises from the scatter in $\Sigma_{\text{SFR}}$, since variations in $\Sigma_{\text{gas}}$ are (by design) narrow for each simulation. The gas depletion time is useful since it is not specific to geometry and can be defined either locally or globally (the area factor cancels out in the definition of $t_{\text{dep}}$). Although values of $t_{\text{dep}}$ in some of our simulations may be somewhat low (see E. Ostriker & C.-G. Kim 2020, in preparation, for discussions on potential causes and missing physics), the scaling may still hold true.

### 5.3. Outflow Velocity, Bernoulli Velocity, and Metal Enrichment

In Figure 11, we present scaling relations for additional wind characteristics, including outflow velocity ($\text{Equation (28)}$), Bernoulli velocity ($\text{Equation (29)}$), and metal enrichment factor ($\text{Equation (26)}$) at $|z| = H$, as a function of $\Sigma_{\text{SFR}}$ with $\tau_{\text{bin}} = 40$ Myr.

Both outflow velocity and Bernoulli velocity scale weakly with $\Sigma_{\text{SFR}}$. The power-law exponent for the $v_{\text{out}}$ versus $\Sigma_{\text{SFR}}$ relation in cool outflows is shallow, $\sim 0.2–0.25$ (depending on where $\tau_{\text{out}}$ is measured). The power-law exponent for the $v_0$ versus $\Sigma_{\text{SFR}}$ relation in hot outflows is even shallower, $\sim 0.07–0.1$ (depending on where $\tau_{\text{out}}$ is measured). These weak scalings seem to be related to the characteristic shell velocity and specific energy of superbubbles driven by SNe at the time of breakout, which are largely insensitive to galactic properties (see Section 6.3).

The hot outflow clearly shows a metal enrichment factor larger than unity, while the cool outflow is only marginally
enriched compared to the bulk of the ISM, and only for high-SFR models. The metal enrichment factor for hot outflows seems to flatten out at low SFRs, so that simple linear regression is not a good description of the behavior. Given the limited number of models, we do not attempt to find a quantitative model from more sophisticated fitting. Instead, we provide simple models shown as the dashed lines in Figures 11(e) and (f) given by

$$\zeta_{\text{cool}} = \begin{cases} 
1.12 (\Sigma_{\text{SFR}}/M_\odot)^{0.05} & \text{if } \Sigma_{\text{SFR}} > 0.1 M_\odot \\
1.0 & \text{otherwise}
\end{cases}$$ (31)
\[
\zeta = \begin{cases} 
(\Sigma_{\text{SFR}}/M_\odot)^{0.15} & \text{if } \Sigma_{\text{SFR}} > 0.1 M_\odot \\
1.5 & \text{otherwise.}
\end{cases}
\] (32)

6. Discussion

6.1. Comparison with Other Simulations

There have been a wide range of local simulations in vertically stratified disks including SN feedback (e.g., Korpi et al. 1999; de Avillez 2000; Joung & Mac Low 2006; Gressel et al. 2008; Joung et al. 2009; Hill et al. 2012; Gent et al. 2013a, 2013b; Hennebelle & Iffrig 2014; Walch et al. 2015; Iffrig & Hennebelle 2017; Peters et al. 2017; Colling et al. 2018; Girichidis et al. 2018b). Notably for present purposes, quantitative analyses of SN-driven outflow properties have been provided in some papers (e.g., de Avillez 2000; Creasey et al. 2013, 2015; Girichidis et al. 2016b; Martizzi et al. 2016; Gatto et al. 2017; Li et al. 2017; Fielding et al. 2018; Kannan et al. 2020), including a few where cosmic rays were part of the physics model (e.g., Girichidis et al. 2016a, 2018a; Simpson et al. 2016). Still, among these studies, only a few have treated SN rates and positions self-consistently with explicit modeling of star formation from self-gravitating collapse, and these have been limited to short-term evolution and considered only a particular galactic condition (e.g., Gatto et al. 2017; Kannan et al. 2020). The present work is the first, to our knowledge, that considers a wide parameter space of local models with different galactic conditions, simulates star formation and feedback self-consistently at high resolution, and follows long-term evolution. As we shall show below, there are interesting similarities and differences between our results on wind scaling and those from other local models.

Comparison with global simulations is of great interest, since these are not subject to some of the limitations of local models. Very recently, cosmological zoom-in simulations have begun to model the disk ISM and star formation without ad hoc subgrid models for wind driving (e.g., Hopkins et al. 2014; Wang et al. 2015; Hopkins et al. 2018b; Marinacci et al. 2019). Quantitative analyses of outflows have been presented, aiming at understanding cosmic baryonic cycles

Figure 11. Scaling relations of characteristic velocities \( v_{\text{out}}, \rho_v \) and metal enrichment factor \( \zeta \) with SFR surface density \( \Sigma_{\text{SFR}} \) at \( |z| = H \). Figures at different heights are available at https://changgoo.github.io/tigress-wind-figureset/figureset.html. The top row ((a) and (b)) is for outflow velocity (Equation (28)), the middle row ((c) and (d)) is for Bernoulli velocity (Equation (29)), and the bottom row ((e) and (f)) is for metal enrichment factor (Equation (26)). All quantities are measured at \( |z| = H \) for cool (left column) and hot (right column) outflows separately. The dashed lines in (e) and (f) denote simple models describing flattening behaviors of \( \zeta \) at low \( \Sigma_{\text{SFR}} \) as in Equations (31) and (32). The simulation results and fitting results are presented as in Figure 8. Characteristic velocities of the hot component are an order of magnitude higher than those of the cool component, although the cool component velocities increase with \( \Sigma_{\text{SFR}} \) slightly more steeply.
(e.g., Muratov et al. 2015, 2017; Christensen et al. 2016; Anglés-Alcázar et al. 2017; Tollet et al. 2019). However, direct comparison of our results on outflow scaling relations with those from zoom-in simulations is beyond the scope of this work because (1) most present scaling of loading factors to global properties, e.g., stellar/halo mass and circular velocity; (2) outflow properties are measured far from their galactic disk origin, and interactions with CGM may have strongly altered the initial outflow properties; (3) outflow analyses mostly do not differentiate by thermal phase (but see Tollet et al. 2019, V. Pandya et al. 2020, in preparation); and (4) although the treatment of the ISM is more realistic than in larger-scale simulations, individual SN feedback is still unresolved, necessitating the adoption of either artificially delayed cooling (Christensen et al. 2016; Tollet et al. 2019) or momentum feedback for most SNe (Muratov et al. 2015, 2017; Anglés-Alcázar et al. 2017). In particular, while the “momentum” feedback approach at mass resolution $\sim 10^3 - 10^5 M_{\odot}$ is able to control star formation, multiphase wind driving from SNe is not resolved.

Idealized global galaxy simulations are a good way to bridge the gap between local and cosmological zoom-in simulations. Currently, only global simulations of dwarf galaxies employ sufficiently high resolution to resolve both star formation and feedback equivalent to this work (e.g., Emerick et al. 2019; Hu 2019). Global simulations of more massive galaxies typically have resolution below what is required to resolve the adiabatic stage of SNR evolution (and therefore to follow hot gas creation), instead adopting “momentum” feedback for most SN events. This is likely why the mass loading of hot gas is lower than that in our simulations; e.g., in the galactic center models of Armillotta et al. (2019) that have similar conditions to our model R2 but mass resolution $2 \times 10^3 M_{\odot}$, the hot gas mass loading factor was $<0.1$, even though for warm gas the fountain-like behavior and mass loading were similar to what we found. In contrast, for the SPH models (MW and Sbc, with mass resolution $500 M_{\odot}$) that have similar conditions to our model R4, Hopkins et al. (2012) found similar mass loading, but in high-velocity escaping rather than moderate-velocity fountain-like warm outflows. This may have been a consequence of the particular implementation of momentum feedback in their particle-based method, which has since been replaced (Hopkins et al. 2018a).

Although an apples-to-apples comparison with existing simulations is not immediately possible, we discuss our results for scaling relations in comparison with work by Creasey et al. (2013, 2015, hereafter C13 and C15, respectively), Martizzi et al. (2016, hereafter M16), and Li et al. (2017, hereafter L17), in which a parameter survey is conducted and scaling relations are presented. We also include results from Smith et al. (2018, hereafter S18), Hu (2019, hereafter H19), and Emerick et al. (2019, hereafter E19). For an in-depth comparison with other simulations (including Girichidis et al. 2016b; Gatto et al. 2017) for solar-neighborhood conditions (model R8), we refer the reader to the discussion in KO18.

Before presenting the results of our comparisons, we begin by summarizing details of the C13, M16, and L17 local simulations, as well as high-resolution global simulations of dwarfs (S18, H19, and E19).

1. C13 conducted non-self-gravitating, unmagnetized local simulations of galactic disks covering a wide range of gas surface density ($2.5 < \Sigma_{\text{gas}} / M_\odot \text{pc}^{-2} < 500$) and gravitational field (parameterized by gas fraction $f_g$ as $g \propto f_g^{1/2}$, which varies from 0.01 to 1). SFRs (and hence SN rates) are prescribed and stay constant over the duration of simulations. They adopt two relations for SFRs: (1) the Kennicutt–Schmidt relation (hereafter C13-KS; Kennicutt 1998),

$$\Sigma_{\text{SFR}} = 2.5 \times 10^{-3} M_\odot \text{Kpc}^{-2} \text{yr}^{-1} \times \left( \frac{\Sigma_{\text{gas}}}{M_\odot \text{pc}^{-2}} \right)^{1.4},$$

and (2) dynamical time prescription (hereafter C13-dyn),

$$\Sigma_{\text{SFR}} = 8.2 \times 10^{-5} M_\odot \text{Kpc}^{-2} \text{yr}^{-1/2} \times \left( \frac{\Sigma_{\text{gas}}}{M_\odot \text{pc}^{-2}} \right)^2.$$

Note that the mean values of $\Sigma_{\text{SFR}}$ in our simulation suite are similar to those of C13-dyn rather than C13-KS, especially at higher surface densities. SNe are placed randomly in the horizontal plane, with the scale height identical to the initial gas scale height. The simulation box is smaller (especially, shorter), $L_x \times L_y \times L_z = 200 \text{pc} \times 200 \text{pc} \times 1 \text{ kpc}$. The fiducial cooling function depends only on density, $n^{1.7} \Lambda$ with constant $\Lambda$, and cuts off at $T = 10^4 \text{ K}$ (some models include a $T$-dependent cooling function). No radiative heating is included.

2. M16 ran non-self-gravitating, unmagnetized local simulations of galactic disks covering $\Sigma_{\text{gas}} = 5, 50$, and $500 M_\odot \text{pc}^{-2}$. The same scaling for SFR surface density is adopted as C13-KS, but the normalization is about a factor of two lower. They have two different SN seeding schemes, but we only compare with their FX models, in which SNe are randomly seeded in space and time within the initial disk scale height. Without a self-consistent treatment of star formation to determine realistic clustering of star formation and hence SNe, their SC models, in which SNe are preferentially seeded near density peaks, result in artificially enhanced cooling of SNe (see also Girichidis et al. 2016b). At their typical resolution of a few parsecs, in their SC models SNe are mostly realized via momentum injection following the prescription of Martizzi et al. (2015), which substantially changes outflow properties (energy loading factor and multiphase structure most significantly). A cubical box with $L = 1 \text{ kpc}$ is adopted. The cooling function depends on temperature but cuts off at $T = 10^4 \text{ K}$. No radiative heating is included.

3. L17 performed non-self-gravitating, unmagnetized local simulations covering $1 < \Sigma_{\text{gas}} / M_\odot \text{pc}^{-2} < 150$. SN rates and distributions are essentially the same as in C13, but additional exploration with independently varying SN scale heights was conducted. The adopted $\Sigma_{\text{SFR}}$ were a bit higher than C13-KS, closer to C13-dyn and our mean $\Sigma_{\text{SFR}}$ (but lower at higher surface densities and higher at lower surface densities). The cooling function depends on temperature and extends to $T \sim 300 \text{ K}$ (Rosen et al. 1993), and a constant photoelectric heating rate was adopted (the effect of the photoelectric heating was explored). The horizontal extent of...
the simulation box is about one-third of ours, but larger (scales with gas surface density) than that of C13.18 4. S18, H19, and E19: the high-resolution global dwarf simulations with which we compare have very low mass (S18: $M_{\text{vir}} = 10^{10} M_\odot$ and $M_{\text{gas}} = 1.8 \times 10^8 M_\odot$; H19: $M_{\text{vir}} = 10^{10} M_\odot$ and $M_{\text{gas}} = 10^7 M_\odot$; and E19: $M_{\text{vir}} = 2.5 \times 10^9 M_\odot$ and $M_{\text{gas}} = 1.8 \times 10^6 M_\odot$). Unlike C13, M16, and L17, all of these simulations have self-consistent star formation and feedback. We note as a caveat in comparison with our results that $\Sigma_{\text{gas}}$ and $\Sigma_{\text{SFR}}$ can vary substantially within a global simulation, and even if one adopts a single value, it will depend on the area. Here we use the scale radius to define the area, but certainly outflows can emerge from locations beyond a scale radius. In the future, more direct comparison with homogeneous definitions would be desirable.

How and where outflow properties are measured matters a great deal, especially for mass loading (see Figure 6). C13 and C15 reported mass, energy, and metal loading factors of the total outflowing gas (without phase separation) by measuring the mean ejected mass, energy, and metals through the vertical boundaries $|z| = 500$ pc. M16 reported the energy loading factor of total outflow gas measured at $1.5 z_{\text{eff}}$, where $z_{\text{eff}}$ is the initial scale height (with slightly different definition from Equation (8)), while they measure mass loading at both $1.5 z_{\text{eff}}$ and 500 pc. L17 reported mass, energy, and metal loading factors for total and hot phases separately by measuring the outflow fluxes averaged over space ($|z| = 1–2.5$ kpc) and time (last 40% of the simulation termination time). H19 and E19 measured outflow fluxes through spherical shells as a function of $r$ rather than $z$. For the purpose of our comparison here, we adopt the compilation of Li & Bryan (2020) at $r = 1$ kpc. To make comparisons as fair as possible, we present plots for total outflow loading factors and match heights as closely as possible given the limitations of reported measurements. We plot C13/C15 results in the panel for $|z| = 500$ pc, M16 results in the panels for $|z| = H$ and 500 pc, and L17, S18, H19, and E19 results in the panel for $|z| = 1$ kpc.

Figure 12 plots $\eta_M$ versus $\Sigma_{\text{gas}}$ measured at (a) $H$, (b) 2H, (c) 500 pc, and (d) 1 kpc, in comparison with the literature results. C13 reported two scaling relations between mass loading factor and gas surface density for two model series, C13-KS and C13-dyn:

$$\eta_M^{\text{C13-KS}} = 13 \pm 10 \left( \frac{\Sigma_{\text{gas}}}{M_\odot \text{ pc}^{-2}} \right)^{-1.15 \pm 0.12} \times f_g^{0.16 \pm 0.14}$$  \hspace{1cm} (35)
Figure 13. Mass loading factor of total outflowing gas as a function of $\Sigma_{\text{SFR}}$, in comparison to other work. For our simulations, mass fluxes are measured at $|z|$ equal to (a) $H$, (b) $2H$, (c) 500 pc, and (d) 1 kpc. The simulation results and fitting results for our models are presented as in Figure 8. Scaling relations reported in C13 are shown in blue (C13-KS; Equation (37)) and green (C13-dyn; Equation (36)) with $f_g = 0.1$ (solid) and 0.5 (dashed). Note that the extent of lines represents the parameter coverage of C13. Magenta stars denote fiducial local models from L17, yellow stars denote local model FX of M16, and cyan symbols show results from global dwarf galaxy models of H19, E19, and S18.

$$\eta_M^{\text{C13-dyn}} = 20 \pm 8 \left(\frac{\Sigma_{\text{gas}}}{M_\odot \, \text{pc}^{-2}}\right)^{-0.82 \pm 0.07} \times f_g^{0.48 \pm 0.08}. \quad (36)$$

In panel (c), we show these two scaling relations for the surface density range consistent with that used in C13. In panel (c), the loading factors we have found are generally higher than in C13 and M16 at a given $\Sigma_{\text{gas}}$. This is mainly because our self-regulated SFR surface densities are higher than their adopted SFR values. From self-regulation, $\Sigma_{\text{SFR}} \propto \Sigma_{\text{gas}} g_s$ is expected, so that the vertical gravity should be taken into account in models assuming a prescribed SFR surface density. Since the boxes in both C13 and M16 are shorter than ours, they generally adopted stronger vertical gravity to confine the gas in the vertical domain. However, their adopted $\Sigma_{\text{SFR}}$ was not adjusted upward corresponding to the expectation from self-regulated star formation at higher $g_s$, and as a result their $\Sigma_{\text{SFR}}$ values are lower than ours at a given $\Sigma_{\text{gas}}$. In addition, stronger gravity would result in higher volume density at a given $\Sigma_{\text{gas}}$, while there would be additional differences in volume filling factors of different gas phases owing to their artificial cooling cutoff. The agreement of C13-dyn with our results is better since the SFR prescription for this model set (implicitly) includes the effect of vertical gravity. The overall better agreement of L17 with our results shown in panel (d) is because the adopted $\Sigma_{\text{SFR}}$ and vertical gravity are more consistent with our simulations (except at the lowest $\Sigma_{\text{gas}}$).

We note that apparent better agreement in panel (a) with M16 is a coincidence, since $z_{\text{eff}}$ is much smaller than the corresponding scale height in our models (see Column (5) in Table 2).

Figure 13 plots the relation between $\eta_M$ and $\Sigma_{\text{SFR}}$ in comparison to the literature results. Using the C13 imposed relation between $\Sigma_{\text{gas}}$ and $\Sigma_{\text{SFR}}$ (Equations (33) and (34)), we can convert Equations (35) and (36) as a function of $\Sigma_{\text{SFR}}$

$$\eta_M^{\text{C13-KS}} = 0.014 \left(\frac{\Sigma_{\text{SFR}}}{M_\odot \, \text{kpc}^{-2} \, \text{yr}^{-1}}\right)^{-0.82} f_g^{0.16}. \quad (37)$$

$$\eta_M^{\text{C13-dyn}} = 0.42 \left(\frac{\Sigma_{\text{SFR}}}{M_\odot \, \text{kpc}^{-2} \, \text{yr}^{-1}}\right)^{-0.41} f_g^{0.07}. \quad (38)$$

In panel (c), we show these two scaling relations for the $\Sigma_{\text{SFR}}$ range consistent with that used in C13. The difference in $\Sigma_{\text{SFR}}$ ranges between the two C13 model series is clearly demonstrated. Again, the C13-dyn results are in fairly good agreement with our results, while C13-KS gives substantially lower mass loading factors. The M16 results in panel (c) also show much lower mass loading factors at a given $\Sigma_{\text{SFR}}$ than our results, as in C13-KS. This implies that $\Sigma_{\text{SFR}}$ is not the only parameter that sets $\eta_M$, but the gravity and/or gas density both matter in setting $\eta_M$.

At 1 kpc, our mass loading factors are again consistent with L17. Overall, we find that the slopes of our mass loading...
relationships are similar at different heights, and these are in good agreement with the literature results in cases where \( \Sigma_{\text{gas}}, g, \) and \( \Sigma_{\text{SFR}} \) are consistent. In addition to local models, in Figure 13(d) we show total mass loading from the global dwarf galaxy simulations of H19, E19, and S18. These results are all fairly consistent with our results. The present agreement may imply that outflow loading factors are independent of the global halo potential. This is encouraging for the development of generalized cosmological subgrid wind launching models from local simulations, although it will be imperative to make further tests and comparisons in other regimes, including more extreme conditions.

We now turn to the energy loading factor. Figure 14 plots our relation between \( \eta_E \) and \( \Sigma_{\text{SFR}} \) for all gas (dominated by the hot medium), in comparison with the literature results. Note that the energy loading factors from C13-dyn are not available, but both C13 and our results suggest that the energy loading factor is insensitive to \( \Sigma_{\text{gas}} \) and \( \Sigma_{\text{SFR}} \). \( \eta_E \sim 0.05-0.5 \) encloses the result reported in C15, which also envelopes our results in panel (c) quite well. L17 obtained a rather narrower range between 0.1 and 0.3, again without strong dependence on \( \Sigma_{\text{SFR}} \). Figure 10 also shows weak dependence of the hot energy loading factors as a function of all galactic parameters we consider. Interestingly, if the energy loading factor is measured at a fixed height, the dependence is even weaker: \( \eta_{E, \text{hot}} \sim 0.1 \) at \( |z| = 500 \) pc and 0.05 at \( |z| = 1 \) kpc. Overall, energy loading factors from our simulations are lower than the fiducial L17 results. The enhanced energy loading factors in L17 are the consequence of larger imposed SN scale heights (\( \sim 150 \) pc) in L17 compared to typical values in our simulations. L17 showed that the energy loading factors increase as the scale height of SNe gets larger since SN explosions in the tenuous disk atmosphere more freely deliver injected energy to the extraplanar region without significant energy loss by cooling (see also the similar tests in Appendix B of KO18). In our simulations, the SN locations are determined by star formation, which in turn depends on the distribution of gas; with “natural” SN positioning with respect to the vertical gas profile, our lower energy loading factors are more consistent with superbubble breakout after shell formation (Kim et al. 2017a).

The energy loading factor can in principle be increased by strongly correlated SNe, since early explosions create a low-density cavity through which energy from subsequent events easily vents with minimal losses (Fielding et al. 2018). Our simulations in fact have highly correlated SNe, with typical cluster particle masses in the range of \( 10^3-10^4 M_\odot \) and with maximum cluster mass up to \( 10^5-10^6 M_\odot \) (higher-mass clusters for inner disk models). However, with long-term evolution and self-consistent inflow/outflow, we find that the energy loading is much reduced compared to the idealized simulations of Fielding et al. (2018). This is because previous or neighboring events can both fill the atmosphere with fountain gas and close off chimneys, both of which render energy venting more difficult (see Figures 1 and 3). Global flow patterns driven by structures like spiral arms and/or bars may potentially reduce inflow/outflow interactions locally; if fountain flows originating in arm regions preferentially fall back in low-density, interarm regions, the energy loading factor may be enhanced in the arm region and reduced in the interarm region, while the
global average stays similar. We are currently analyzing the outflow properties from local simulations with spiral arms utilizing the TIGRESS framework (Kim et al. 2020). Higher energy loading factors in global dwarf simulations (as, e.g., shown from H19, E19, and S18 in Figure 10) may also be related to the global geometry, but caution needs to be taken since the cooling rates in dwarf simulations are generally lower owing to the lower metallicity.

In Figure 15, we compare the correlation between our energy and SN-origin metal loading factors, in comparison with the literature results (see also Li & Bryan 2020). Since SNe drive outflows, the fluxes of energy and SN-origin metals have a common origin and are expected to correlate with each other, and C15 previously identified a tight correlation. Note that if we use total metal fluxes, this correlation gets weaker (almost disappears). In our simulations, even with $Z_{SN} = 10Z_{ISM,0}$, more mass comes from the ISM than from SNe, so the metal mass flux is dominated by the ISM-origin metals (see Figure 5). We consider both cumulative and instantaneous SN-origin metal fluxes. The former is for metals injected by SNe over the entire simulation duration and directly measurable from the metal tracer field employed in the simulation. The latter is for metals injected by recent SN events and obtained from Equation (25) (see Section 4.2). Note that although reported metal loading factors in other simulations technically correspond to cumulative SN-origin metals, their metal loadings can be interpreted as instantaneous ones since C13 and L17 run for a much shorter time than we do, and H19 and E19 consider low-metallicity dwarfs.

Figure 15 plots energy loading factors as a function of cumulative (a) and instantaneous (b) SN-origin metal loading factors measured at $|z| = 1$ kpc. The top row ((a) and (b)) is for the total outflow, and the bottom row ((c) and (d)) is for the hot outflow. The dotted line is $\eta_E = \eta_Z$, and the dashed line is $\eta_E = 0.4\eta_Z$ as suggested by C15. In contrast to other scaling relations, we use all points from time evolution for fitting since temporal correlations within a model between the two loading factors are in fact physically meaningful; there is no temporal offset.

| Figure 15. Correlation between energy loading factor and SN-origin metal loading factor at $|z| = 1$ kpc, with comparison to other work. Figures at different heights are available at https://changgoo.github.io/tigress-wind-figureset/figureset.html. Left (a) and (c) and right (b) and (d) columns are for cumulative and instantaneous measures of SN-origin metal loading factor obtained by Equation (25) with initial and instantaneous ISM metallicity, respectively. Top (a) and (b) and bottom (c) and (d) rows are for total and hot outflows, respectively. The reference lines for $\eta_E = \eta_Z$ (dotted) and $\eta_E = 0.4\eta_Z$ (dashed; C15) are also shown. The orange region covers the result of C15. The fitting results for our models are presented as in Figure 8. |
between two fluxes, and loading factors use the same denominator (up to a constant factor).

Without radiative energy loss, the energy loading factor would be equal to the (instantaneous, SN-origin) metal loading factor. As energy is lost by radiative cooling, \( \eta_E < 1 \) and \( \eta_E < \eta_Z \) (C15). In addition, due to “recycled” metals through fountain flows, the cumulative SN-origin metal flux is larger than the instantaneous one. However, energy in fountain flows is radiated away and not “recycled,” so that the ratio \( \eta_E/\eta_{SN}^{\text{cum}} \approx 0.09 \) is smaller than \( \eta_E/\eta_{SN} \approx 0.3 \), as shown in Figures 15(a) and (b). The relation reported in C15 for the total outflowing gas, \( \eta_E = 0.4\eta_Z \), is quite close to our result for the instantaneous metal flux measurement.

As expected, the correlation gets tighter when only hot outflows are considered (i.e. and (d)) since cooling is minimal in hot outflows. The energy-to-metal loading factor ratio is also increased with the instantaneous metal loading factor. The slope in hot outflows is steeper than unity, implying less efficient cooling when there is more efficient loading of SN-origin metals in hot outflows. In other words, successful breakouts due to clustered SN events (indicated by high SN-origin metal loading factor) load SN energy to outflows more efficiently (Fielding et al. 2018). Using the fitting result in Figure 15(d),

\[
\eta_{E,\text{hot}} = 0.81 \eta_{Z,\text{hot}}^{1.15},
\]

and Equation (29), we obtain

\[
f_{E,\text{hot}} = \left( \frac{2E_{SN}}{m_{\*}} \right)^{1/2} \left( \frac{\eta_{E,\text{hot}}}{\eta_{M,\text{hot}}} \right)^{1/2}
= 2.9 \times 10^3 \text{ km s}^{-1} \eta_{SN,\text{hot}}^{0.58} \eta_{M,\text{hot}}^{0.08}.
\]

This says that the specific energy in hot outflows is most sensitive to the fraction of genuine SN material in outflows,

\[
f_{SN,\text{hot}} = \frac{Z_{\text{hot}} - Z_{\text{ISM}}}{Z_{\text{SN}} - Z_{\text{ISM}}},
\]

which varies from event to event. To enhance \( f_{SN,\text{hot}} \) and hence \( \eta_{E,\text{hot}} \) on average, SN feedback needs to occur either preferentially outside the main gas disk (L17) or inside a region in which a vertical cavity has been opened. The latter case is not easily realized in our simulations but may be possible in central starbursts.

6.2. Comparison with Observations

Observations of galactic outflows (winds) are challenging because the outflow is much more tenuous than the underlying galactic disk, so that both emission and absorption lines are weaker. At the same time, outflows possess complex, multi-phase structure, demanding high-sensitivity observations of many gas tracers to quantify the mass (total and metal), momentum, and energy budget of the outflow. Currently, direct observational constraints for outflow characteristics and their scaling relations with galactic properties are neither strong nor comprehensive (see Rupke 2018, for a review).

Optical and UV absorption-line surveys provide the largest body of data to study correlations between the outflow characteristics and galaxy properties (Martin 2005; Rupke et al. 2005; Arribas et al. 2014; Chisholm et al. 2015; Heckman et al. 2015; Cazzoli et al. 2016; Heckman & Borthakur 2016). From line profiles, it is relatively straightforward to derive the characteristic velocity of the outflow (modulo different definitions adopted in different studies). A shallow, positive correlation between outflow velocity and SFR is consistently observed in both neutral and ionized outflows: \( \eta_{out} \propto M_\*^{0.15-0.35} \). Heckman & Borthakur (2016) presented a similar correlation between outflow velocity and SFR surface density, \( \eta_{out} \propto \Sigma_{\text{SFR}}^{0.34} \) (essentially the same correlation is seen in stacking analysis of galaxies at \( z \approx 2 \) by Davies et al. 2019), while Chisholm et al. (2015) did not find a convincing correlation of \( \eta_{out} \) with \( \Sigma_{\text{SFR}} \). Consistent with the observations, we find weak scalings, approximately \( \eta_{out} \propto \Sigma_{\text{SFR}}^{0.2} \), for both hot and cool gas, but an order of magnitude higher velocity for the former (Figure 11(a)). We note that these observations treat galaxies as a whole and are therefore not directly equivalent to our scaling relations (which would require observational resolution of \( \lesssim \) kpc and sufficient sensitivity to detect individual disk “patches”). Also, the range of \( \Sigma_{\text{SFR}} \) in observations described above is generally on the high side, \( \Sigma_{\text{SFR}} > 0.1 M_\* \text{kpc}^{-2} \text{yr}^{-1} \), which only marginally overlaps with our parameter space.

The mass loading factor is a more difficult quantity to measure empirically. In estimating the mass loading from observed interstellar absorption lines, many assumptions are involved, including the covering area of the outflow (a combination of the opening angle, characteristic radius, and covering fraction of the outflow), the column density conversion from a specific species to total hydrogen, and the characteristic velocity (e.g., Rupke et al. 2005). The reported mass loading factors from observations of dwarf starbursts and LIRGs/ULIRGs are in the range \( \eta_{M} \sim 0.1-10 \) and have found either negative correlation (e.g., Heckman et al. 2015; Chisholm et al. 2017) or no correlation (e.g., Martin 1999; McQuinn et al. 2019) with galaxy mass (or circular velocity). Although the full galaxy mass range in these studies is \( \log M_\* \sim 7-11 \), the low-mass galaxy samples (at \( \log M_\* \sim 7-8 \)) used in the study that found negative correlation are more extreme starbursts than those in the study that reported no correlation (see McQuinn et al. 2019). Arribas et al. (2014) observed local LIRGs and ULIRGs (\( \log M_\* \sim 9.5-11 \)) with integral field spectroscopy and obtained \( \eta_{M} \propto M_\*^{-0.43} \), similar to Chisholm et al. (2017). A direct comparison with our results is not possible, since our work measures outflow rates and galactic properties locally, in contrast to the global outflow rates and galaxy mass reported in observations. Still, it is encouraging that the observed estimates of \( \eta_{M} \) are similar to what we find (Figure 13) at \( \Sigma_{\text{SFR}} \sim 0.1-1 \), which overlaps with the observed range for these samples.

Interestingly, Arribas et al. (2014) reported a positive correlation between \( \eta_{M} \) and \( \Sigma_{\text{SFR}} \) with a log–log slope of 0.17, which is apparently in tension with our results (see Figure 9, with slopes \( \sim 0.5 \) for cool gas) and those from other numerical simulations, which all show negative scaling for \( \eta_{M} \) versus \( \Sigma_{\text{SFR}} \) (Figure 13). However, in Arribas et al. (2014) \( \Sigma_{\text{SFR}} \sim 0.1-100 M_\* \text{kpc}^{-2} \text{yr}^{-1} \), which only marginally overlaps with the high end of our \( \Sigma_{\text{SFR}} \) range. Furthermore, the scatter in their mass loading factor is large, and the significance
of the fit is not high (Figure 14 of Arribas et al. 2014). Nevertheless, there is overall agreement in the range of mass loading factor, $\eta_{fl} \sim 0.1$–1. Our results also suggest an intriguing possibility of a weakened correlation between $\eta_{fl}$ and $\Sigma_{SFR}$ at high $\Sigma_{SFR}$, where hot outflows begin to dominate the total mass (Figures 8(a) and (b)).

In the future, spatially resolved outflow observations utilizing sensitive integral field unit observations offer the promise of enabling direct comparison with the kind of local scaling relations reported here. With future computational advances, it will also be possible to run global simulations with the current resolution and physics of our local simulations, to connect with observed global relationships.

6.3. Physical Interpretation of Scaling Relations

Multiphase outflow launching in our simulation suite is an outcome of intricate interactions between SN feedback and ISM dynamics, with complexity that precludes a purely analytic theory that can explain our quantitative findings. Nevertheless, we are able to obtain insight into the physics behind the emergent scaling relations we have found using a simple theoretical model of superbubble evolution and breakout. Given the simple assumptions we adopt (e.g., uniform-background medium and spherical symmetry), we will mainly focus on parameter dependence rather than coefficients.

Weaver et al. (1977) developed an analytic theory for the evolution of stellar-wind-blown bubbles, and essentially the same theory has subsequently been applied to superbubbles driven by clustered SNe (McCray & Kafatos 1987; Mac Low & McCray 1988; El-Badry et al. 2019). In the model, the evolution after the radiative shell formation is characterized by the energy injection rate $\dot{E}_{in}$ and the ambient medium density $\rho_0$. In Weaver et al. (1977), the injected energy is shared among kinetic energy of the cooled shell $\dot{E}_{kin,cool}$, thermal energy of the hot interior $\dot{E}_{th,hot}$, and radiative energy losses in the forward shocks $\dot{E}_{shock-cooling}$. The classical theory predicts $\dot{E}_{kin,cool} = (15/77)\dot{E}_{in}$, $\dot{E}_{th,hot} = (5/11)\dot{E}_{in}$, and $\dot{E}_{shock-cooling} = (27/77)\dot{E}_{in}$. Since $\dot{E}_{in} \propto$ SFR, to zeroth order this explains why energy loading factors of both cool and hot outflows are nearly constant with SFR (see Figures 8(e) and (f)).

The classical theory neglects cooling at the interface between the hot interior and cool, dense shell, while in reality the interface cooling $\dot{E}_{interface-cooling}$ is crucial for understanding the energy budgets in superbubbles (e.g., Kim et al. 2017a; Fielding et al. 2018; Gentry et al. 2019). Mixing layers between hot and cool gas are mediated by both (M)HD instabilities and radiative cooling, best explored with very high resolution simulations (e.g., Fielding et al. 2020a). In the current simulations, the existence of intermediate-temperature phase in outflows demonstrates that cooling in the mixing layers plays a role in reducing the injected energy.

For present purposes, we employ a model used in 1D simulations of El-Badry et al. (2019), in which interface mixing is parameterized via a diffusion coefficient $\lambda \nu$. The resulting interface cooling rate is $\dot{E}_{interface-cooling} = \theta \dot{E}_{in}$, with $\theta$ depending on $\lambda \nu$ and ambient density $\rho_0$ as $\theta/(1 - \theta) \propto (\lambda \nu)^{1/2} \rho_0^{1/2}$. Inclusion of the interface cooling reduces $\dot{E}_{in}$ to $(1 - \theta)\dot{E}_{in}$, and with less power the bubble expands less rapidly. This results in $\dot{E}_{th,hot} = (5/11)(1 - \theta)\dot{E}_{in}$ and $\dot{E}_{shock-cooling} = (27/77)(1 - \theta)\dot{E}_{in}$, so that a constant $\theta$ would still imply energy loading of hot and cool outflows that are independent of SFR. In reality, the ambient medium in the real ISM (and in the current simulations) is highly inhomogeneous and vertically stratified, and the diffusion coefficient representing details of mixing layer varies, so that $\theta$ is not constant. The weak scaling between $\eta_{fl}$ and $\Sigma_{SFR}$ (Figures 8(e) and (f)) presumably arises from weak dependencies in the averages over these variations.

Since a superbubble’s interior temperature depends very weakly on the ambient medium density ($T_{hot} \propto \rho_0^{-2/35}$ for conduction-mediated evaporation from Weaver et al. 1977; El-Badry et al. 2019; and $T_{hot}$ is also insensitive to $\rho_0$ from simulations of expansion in an inhomogeneous medium without conduction from Kim et al. 2017a), the constant mass loading factor of hot outflows (Figure 8(b)) is easily understood from $\eta_{fl,hot} \sim \eta_{E,hot}/T_{hot}$ with weakly varying $T_{hot}$.

For the mass loading factor of cool outflows, $\eta_{fl,cool} \sim \eta_{E,cool}/T_{cool}^{2}$ we need to understand what determines the characteristic outflow velocity of the cool phase. To this end, we seek a scaling relation of the cooled shell velocity when a superbubble breaks out of the disk (roughly $R \sim H$). Applying the theory of El-Badry et al. (2019), the bubble radius follows

$$R(t) = 83 \text{ pc}(1 - \theta)^{1/5} \left(\frac{\dot{E}_{in}}{10^{46} \text{ erg yr}^{-1}}\right)^{1/5} \times \left(\frac{n_0}{\text{cm}^{-3}}\right)^{-1/5} \left(\frac{t}{\text{Myr}}\right)^{3/5},$$

and the shell velocity is

$$v_{sh}(t) = 49 \text{ km s}^{-1}(1 - \theta)^{1/5} \left(\frac{\dot{E}_{in}}{10^{46} \text{ erg yr}^{-1}}\right)^{1/5} \times \left(\frac{n_0}{\text{cm}^{-3}}\right)^{-1/5} \left(\frac{t}{\text{Myr}}\right)^{-2/5}.\quad(43)$$

The time at which the bubble radius reaches the disk scale height is

$$t_H \equiv 8.5 \text{ Myr}(1 - \theta)^{-1/3} \left(\frac{\dot{E}_{in}}{10^{46} \text{ erg yr}^{-1}}\right)^{-1/3} \times \left(\frac{n_0}{\text{cm}^{-3}}\right)^{1/3} \left(\frac{H}{300 \text{ pc}}\right)^{5/3},$$

so that

$$v_{sh}(t_H) = 21 \text{ km s}^{-1}(1 - \theta)^{1/3} \left(\frac{\dot{E}_{in}}{10^{46} \text{ erg yr}^{-1}}\right)^{1/3} \times \left(\frac{n_0}{\text{cm}^{-3}}\right)^{-1/3} \left(\frac{H}{300 \text{ pc}}\right)^{-2/3}.\quad(45)$$
Assuming that all SNe that explode within an area $\pi R^2$ contribute to superbubble breakout at $R = H$,

$$E_{\text{in}}(<H) = \pi H^2 E_{\text{SN}} \frac{\Sigma_{\text{SFR}}}{m_\odot} = 3 \times 10^{47} \text{ erg yr}^{-1} \times \left( \frac{\Sigma_{\text{SFR}}}{0.1 M_\odot \text{ kpc}^{-2} \text{ yr}^{-1}} \right) \left( \frac{H}{300 \text{ pc}} \right)^2. \quad (46)$$

We obtain

$$v_{\text{sh}}(t_H) = 63 \text{ km s}^{-1} (1 - \theta)^{1/3} \left( \frac{\Sigma_{\text{SFR}}}{0.1 M_\odot \text{ kpc}^{-2} \text{ yr}^{-1}} \right)^{1/3} \times \left( \frac{n_0}{\text{ cm}^{-3}} \right)^{-1/3}. \quad (47)$$

Since SFRs in our simulations agree well with the pressure-regulated, feedback-modulated star formation theory (E. Ostriker & C.-G. Kim 2020, in preparation; see also Ostriker et al. 2010; Kim et al. 2011; Ostriker & Shetty 2011), we may use the relationships $W = P_{\text{mid}} = \Upsilon \Sigma_{\text{SFR}}$, where $\Upsilon$ is the total feedback yield$^{20}$ (allowing for thermal and magnetic as well as turbulent terms). We assume the characteristic ambient medium density to be the midplane density $\rho_{\text{mid}}$, which is

$$\rho_{\text{mid}} = \frac{P_{\text{mid}}}{\sigma_{z,\text{eff}} \Sigma_{\text{SFR}}} = \frac{\Upsilon \Sigma_{\text{SFR}}}{\sigma_{z,\text{eff}}}, \quad (48)$$

or

$$n_{\text{mid}} = 1.7 \text{ cm}^{-3} \left( \frac{\Upsilon}{10^3 \text{ km s}^{-1}} \right) \left( \frac{\Sigma_{\text{SFR}}}{0.1 M_\odot \text{ kpc}^{-2} \text{ yr}^{-1}} \right) \times \left( \frac{\sigma_{z,\text{eff}}}{40 \text{ km s}^{-1}} \right)^{-2}. \quad (49)$$

We then finally obtain

$$v_{\text{sh}}(t_H) = 52 \text{ km s}^{-1} (1 - \theta)^{1/3} \left( \frac{\Upsilon}{10^3 \text{ km s}^{-1}} \right)^{-1/3} \times \left( \frac{\sigma_{z,\text{eff}}}{40 \text{ km s}^{-1}} \right)^{2/3}, \quad (50)$$

with no explicit dependence of $v_{\text{sh}}(t_H)$ on $\Sigma_{\text{SFR}}$. Note that in previous work seeking a physical interpretation of the observed weak scaling between the outflow velocity and $\Sigma_{\text{SFR}}$, the empirical Kennicutt–Schmidt relation $\Sigma_{\text{SFR}} \propto \Sigma_\ast^{1.4}$ (Kennicutt 1998) was instead adopted to get $v_{\text{sh}}(t_H) \propto \Sigma_{\text{SFR}}^{0.1 \pm 0.17}$. Modulo a hidden dependence in $(1 - \theta)$, this explains the weak, positive scaling $\sigma_{\text{out,cool}} \propto \Sigma_{\text{SFR}}^{0.23}$ and hence $\eta_{\text{sh},\text{cool}} \propto \Sigma_{\text{SFR}}^{0.23}$, similar to the results shown in Figures 11(a) and 8(a).

We emphasize that $\eta_{\text{sh},\text{cool}}$ is a characteristic velocity from a rather wide distribution of $v_{\text{out}}$ rather than a single “shell” velocity $v_{\text{sh}}(t_H)$ as in the above simple theory. Even in idealized simulations of multiple SNe in an inhomogeneous medium (Kim et al. 2017a), the distribution of expanding velocities is broad, while the characteristic “knee” in the velocity distribution increases with the energy injection rate (parameterized by an interval between SNe) but is insensitive to density.

### 6.4. TIGRESS Outflow Models in Context

The methods used in this work have clear pros and cons in the context of galactic wind research. Here we review the advantages and also discuss limitations of our methodology.

With the uniformly high resolution of our simulations (2–8 pc; higher resolution for denser condition), the outflow characteristics studied in this work arise not from ad hoc assumptions but from resolved key physical processes at every relevant step:

1. **Star formation.**—Star formation occurs in gravitationally collapsing objects at high density and pressure that is distinct from the ambient ISM (e.g., Mao et al. 2020).
2. **SN injection.**—Self-regulated SFRs and a population synthesis model applied to star cluster particles provide SN rates and positions that have realistic spacetime correlations with respect to each other and the distribution of ISM gas.
3. **Superbubble evolution.**—The Sedov–Taylor stage of SNR evolution is resolved for more than 90% of individual SNe, directly capturing hot gas creation and momentum injection.
4. **Multiphase outflow evolution.**—The evolution of low-density outflows in extraplanar regions is followed using the same spatial and time resolution as the higher-density ISM near the midplane (Vijayan et al. 2020), without degrading the resolution as in (semi-)Lagrangian or adaptive mesh refinement schemes.
5. **Long-term evolution.**—Each model is run at least up to $1.5 t_{\text{dyn}}$, covering a few star formation feedback wind launching outflow/inflow cycles.

The main caveats arise from the local approximation (adopted to achieve uniformly high resolution) and missing physics (adopted to enable long-term evolution and a survey of parameters):

1. **Missing global geometry.**—Outflow evolution to scales large compared to the launch region cannot be captured in local models. Without streamline opening, hot winds do not reach their asymptotic velocity (e.g., Chevalier & Clegg 1985; Fielding et al. 2017; Smith et al. 2018), and fountain flows that travel large radial distances cannot be captured.
2. **Missing radial and cosmic accretion.**—Our simulation adopts outflow boundary conditions in the vertical direction and shearing-periodic boundary conditions in the horizontal directions. There are therefore no sources of new gas to replace gas lost to star formation or winds. It is worth emphasizing that the galactic-scale impact of outflows would be determined not solely by wind launching properties characterized in this paper but also by interaction with the CGM, which is in part shaped by cosmic flows that cannot be modeled in local simulations (Fielding et al. 2020b). The relevant processes include cosmic accretion, gas flows driven by galaxy mergers, and outflows from satellite galaxies.
3. **Missing early feedback.**—We only include the two dominant channels of stellar feedback, SNe and radiative
heating of warm–cold gas. It has previously been argued that dynamics driven by “early feedback” in the form of radiation pressure, massive-star winds, and photoionization is needed to reduce densities and make SNe effective (Gatto et al. 2017; Peters et al. 2017; Kannan et al. 2020). In fact, the natural clustering of SNe in our simulations means that we fully resolve radiative SNR evolution >90% of the time. However, in environments where the free-fall times in dense clouds are short, the lack of early feedback means that star clusters may significantly grow in the ~3–4 Myr before the onset of the first SN; this may be responsible for unrealistically high SFRs in our models R2 and R4. For lower-density environments, SNe effectively disperse their parent clouds without excessive star formation. In (short-term) simulations with conditions similar to model R8, Gatto et al. (2017) found similar galactic outflow fluxes for models with and without stellar winds, while Kannan et al. (2020) found similar outflow fluxes for models with and without radiation pressure.

4. Other missing physics.—Thermal conduction and cosmic rays are two major missing physical processes that may have potentially significant impact on our results. Thermal conduction can load more hot gas during the superbubble evolution (e.g., El-Badry et al. 2019). Since superbubbles in our simulations expand in a highly inhomogeneous, turbulent ISM, there is a high level of mixing that can transfer gas between warm and hot phases (see also Schneider et al. 2020, for evidence of this). It remains unclear whether fully realistic simulations that also include thermal conduction (which must be anisotropic to allow for the magnetic field) alter mass loading of hot outflows significantly.

Cosmic rays are mainly accelerated in SN shocks and provide a nonthermal pressure force with relatively low radiative losses. Cosmic rays advect with the gas and also diffuse along the magnetic field, with flux limited by the Alfvén speed. Although there are large uncertainties in diffusion coefficients and numerical difficulties in modeling cosmic-ray transport, cosmic-ray pressure gradients may be substantial and play a key role in driving cooler, smoother, and slower galactic winds (e.g., Simpson et al. 2016; Girichidis et al. 2018a; Mao & Ostriker 2018).

7. Summary

This work quantifies characteristics of multiphase outflows emerging from self-consistent, high-resolution simulations of the star-forming ISM. Our suite of MHD simulations consists of seven models covering a range of galactic conditions that appear within normal star-forming galaxies like the Milky Way. Each model represents a local, roughly kiloparsec-scale region within a galactic disk. The ISM in each simulation is explicitly modeled by solving ideal MHD equations, including the effect of galactic differential rotation, gas self-gravity, external gravity from the stellar disk and dark matter halo, optically thin cooling from 10 to 10^9 K, photoelectric heating onto small grains by FUV radiation, and energy and momentum input from SNe. Gas collapses to make star cluster particles, which produce in situ and runaway SNe. The TIGRESS framework (see KO17 for numerical details) allows us to follow long-term evolution (more than an orbit time, at least a few feedback cycles after the initial transient) of the star-forming ISM, with self-regulated SFRs and ISM properties. Self-regulation cycles of star formation and feedback modulate outflows and inflows self-consistently (Figure 3). Galactic winds emanating from superbubble breakout possess multiphase structure with distinct characteristics (see KO18 and Vijayan et al. 2020 for in-depth analysis of the solar-neighborhood model R8). We measure fluxes of mass (total and metal), momentum, and energy of each thermal phase of the outflowing gas at four different locations: |z| = H, 2H, 500 pc, and 1 kpc; results are given in Section 4. We present scaling relations for wind loading factors, characteristic velocities, and metal properties as a function of a variety of local galactic properties. These scaling relations are reported separately for cool and hot phases (Section 5), and we also compare scalings of total loading with results from other recent simulations (Section 6.1) and observations (Section 6.2). We provide a physical interpretation of scalings based on a simple theoretical model of superbubble breakout (Section 6.3). We provide full information from our outflow analyses at doi:10.5281/zenodo.3872049, which we hope can serve as a benchmark for upcoming theoretical and observational studies.

Our key findings for galactic outflows are as follows:

1. Overall evolution.—Star formation, SN feedback, and wind driving are all self-regulated and show clear cyclic behavior (Section 3). In low surface density models, the characteristic timescale for vertical oscillation (π√gρν)_1/2 is longer than the feedback timescale (or star cluster evolution timescale ~40 Myr), leading to a well-defined cyclic behavior for star formation and outflow fluxes governed by vertical oscillation. In high surface density models, in contrast, the natural vertical oscillation period is shorter than the duration of feedback from a burst, so that returning flows interfere with gas being launched by a burst. In these cases, evolution is more chaotic and no clear correspondence between midplane starbursts and outflows above the disk exists. We thus construct time-averaged outflow characteristics over a few feedback cycles (0.5 < t/t_{orb} < 1.5) to quantify the overall behavior, rather than individual bursts. This is especially important in the measurement of “loading factors,” for which a mismatch between time-dependent outflow fluxes and offset time-dependent reference fluxes (set by SN/SFRs) can produce quite misleading instantaneous measurements for loading.

2. Emergent multiphase outflow ranges.—For the ranges Σ_{gas} ~ 1–100 M_{⊙} pc^-2 and Σ_{SN}/(2ζ_z + ρ_{tm}) = 0.005–1 M_{⊙} pc^-3 that are inputs to our simulations, the ranges of self-consistently regulated properties of the star-forming ISM disk are Σ_{SFR} ~ 10^{-4}–1 M_{⊙} kpc^-2 yr^-1, P_{mid} ~ W ~ 10^3–10^{10} erg cm^-3 K, n_{mid} ~ 0.05–50 cm^-3, and t_{dep} ~ 10^5–10^6 Myr. From fluxes measured at |z| = H, the emergent loading factors of mass, momentum, and energy are η_{m} ~ 0.5–50, η_{p} ~ 0.04–0.7, η_{E} ~ 0.005–0.02, η_{SN} ~ 0.1 for cool outflows (T < 2 × 10^4 K) and η_{m} ~ 0.1–0.7, η_{p} ~ 0.07–0.12, η_{E} ~ 0.05–0.25, η_{SN} ~ 0.1–0.3 for hot outflows (T > 5 × 10^4 K). The intermediate phase (2 × 10^4 K < T < 5 × 10^5 K) is subdominant for all loading factors.
We note that at fixed height the hot outflow energy loading factor is essentially constant across simulations (e.g., Figures 14(c) and (d)), \( \eta_{E,\text{hot}} \approx 0.1 \) for \( |z| = 500 \text{ pc} \) and \( \eta_{E,\text{hot}} = 0.04 \) at 1 kpc. Similarly to the energy loading factor, at fixed heights \( |z| = 500 \text{ pc} \) and 1 kpc, the instantaneous SN-origin metal loading factor is more or less constant, \( \eta_{\text{SN},Z,\text{hot}} = 0.16 \) and 0.066, respectively. Figures and the data at different heights are available at doi:10.5281/zenodo.3872049.

3. Scaling of loading factors.—We find that mass is primarily carried by cool outflows and energy is primarily carried by hot outflows, with the following scaling relations for loading factors at \( |z| = H \):

\[
\log \eta_{M,\text{cool}} = -0.44^{+0.08}_{-0.08} \log \left( \frac{\Sigma_{\text{SFR},40}}{M_\odot \text{ kpc}^{-2} \text{ yr}^{-1}} \right) - 0.07^{+0.16}_{-0.15} \pm 0.27 \quad \text{Figure 9(b)} \tag{51}
\]

\[
= -0.54^{+0.08}_{-0.08} \log \left( \frac{W/k_B}{\text{cm}^{-3} \text{ K}} \right) + 3.23^{+0.39}_{-0.42} \pm 0.23 \quad \text{Figure 9(g)} \tag{52}
\]

\[
= 0.70^{+0.09}_{-0.10} \log \left( \frac{t_{\text{dep},40}}{\text{Myr}} \right) - 1.44^{+0.32}_{-0.29} \pm 0.23 \quad \text{Figure 9(h)} \tag{53}
\]

\[
\log \eta_{E,\text{hot}} = 0.14^{+0.08}_{-0.08} \log \left( \frac{\Sigma_{\text{SFR},40}}{M_\odot \text{ kpc}^{-2} \text{ yr}^{-1}} \right) - 0.70^{+0.12}_{-0.14} \pm 0.22 \quad \text{Figure 9(h)} \tag{54}
\]

\[
= 0.17^{+0.09}_{-0.09} \log \left( \frac{W/k_B}{\text{cm}^{-3} \text{ K}} \right) - 1.73^{+0.49}_{-0.47} \pm 0.21 \quad \text{Figure 10(g)} \tag{55}
\]

\[
= -0.25^{+0.11}_{-0.10} \log \left( \frac{t_{\text{dep},40}}{\text{Myr}} \right) - 0.27^{+0.30}_{-0.32} \pm 0.20 \quad \text{Figure 10(h).} \tag{56}
\]

The variation of mass loading factors with galaxy properties is strong in cool outflows and weak in hot outflows. In fact, \( all \) loading factors of hot outflows only vary by a factor of 2–3 (right column of Figure 8), while galactic properties like \( \Sigma_{\text{SFR}} \) vary by more than 4 orders of magnitude.

For cool gas, the momentum loading also varies significantly across galaxy environments, while energy and metal loading do not (left column of Figure 8). We find overall a similar level of correlations between loading factors and all local galactic properties we consider except \( \Sigma_{\text{gas}} \) (Figure 9, Figure 10). This is in part because the “derived” galactic properties \( (\Sigma_{\text{SFR}}, P_{\text{mid}}, W, n_{\text{mid}}, \text{ and } t_{\text{dep}}) \) are self-regulated and connected with each other, and in part because our parameter choice assumes an implicit correlation between gas \( (\Sigma_{\text{gas}}) \) and gravity \( (\Sigma_{\phi}/\Sigma_{*}) \) and \( \rho_{\text{dm}} \) parameters (see Appendix C). Subsequent work exploring a wider parameter space would be needed to cover conditions in nearby observable targets, including dwarf starbursts and LIRGs/ULIRGs, and the full range of conditions that are relevant to theoretical galaxy formation models (Motwani et al. 2020).

We emphasize that the large mass loading of outflows at low SFR does not imply a massive cool wind in high-mass galaxies because the cool gas outflow velocities are low. Instead, at low SFR there is a heavily loaded cool fountain.

4. Characteristic velocities.—We define two characteristic velocities, an outflow velocity \( \tau_{\text{out}} \) (Equation (28)) and a Bernoulli velocity \( \tau_{\text{g}} \) (Equation (29)). Since we include all gas that has positive outward velocity in computing outflow fluxes, the low \( \tau_{\text{out}} \sim 10^{-1} \text{ km s}^{-1} \) values we find for cool-phase outflows imply that in high-mass galaxies a large fraction of the gas will fall back as fountains, as indeed the simulations show. For cool outflows, \( \tau_{\text{g}} \sim 20–140 \text{ km s}^{-1} \) is dominated by the kinetic term and is not much larger than \( \tau_{\text{out}} \). For hot outflows, \( \tau_{\text{g}} \sim 400–1400 \text{ km s}^{-1} \) is dominated by the thermal term and is large enough that hot gas would escape far into halos for most galaxies. We find generally weak scaling of the characteristic velocities with galactic properties.

The velocity scaling relations at \( |z| = H \) obtained in this work are

\[
\log \left( \frac{\tau_{\text{out,cool}}}{\text{km s}^{-1}} \right) = 0.23^{+0.04}_{-0.03} \log \left( \frac{\Sigma_{\text{SFR},40}}{M_\odot \text{ kpc}^{-2} \text{ yr}^{-1}} \right) + 1.78^{+0.07}_{-0.07} \pm 0.14 \quad \text{Figure 11(a)} \tag{57}
\]

\[
= 0.27^{+0.03}_{-0.03} \log \left( \frac{W/k_B}{\text{cm}^{-3} \text{ K}} \right) + 0.10^{+0.17}_{-0.10} \pm 0.10 \quad \text{Table 5} \tag{58}
\]

\[
= -0.34^{+0.03}_{-0.04} \log \left( \frac{t_{\text{dep},40}}{\text{Myr}} \right) + 2.46^{+0.11}_{-0.11} \pm 0.08 \quad \text{Table 5} \tag{59}
\]

\[
\log \left( \frac{\tau_{\text{g,hot}}}{\text{km s}^{-1}} \right) = 0.11^{+0.04}_{-0.04} \log \left( \frac{\Sigma_{\text{SFR},40}}{M_\odot \text{ kpc}^{-2} \text{ yr}^{-1}} \right) + 3.04^{+0.08}_{-0.08} \pm 0.16 \quad \text{Figure 11(d)} \tag{60}
\]

\[
= 0.13^{+0.05}_{-0.05} \log \left( \frac{W/k_B}{\text{cm}^{-3} \text{ K}} \right) + 2.25^{+0.23}_{-0.21} \pm 0.14 \quad \text{Table 5} \tag{61}
\]

\[
= -0.17^{+0.06}_{-0.05} \log \left( \frac{t_{\text{dep},40}}{\text{Myr}} \right) + 3.37^{+0.16}_{-0.18} \pm 0.14 \quad \text{Table 5}. \tag{62}
\]

5. Metals.—Metals in outflows originate from both the ISM and SN. Recent SN-origin material in hot outflows amounts to typically 5%–20% of the mass and 30%–60% of the metal mass (these fractions generally increase with \( \Sigma_{\text{SFR}} \)).

The instantaneous SN-origin metal loading factor scales very weakly with all galactic properties, e.g., at \( |z| = H \),

\[
\log \eta_{\text{SN},Z,\text{hot}} = 0.11^{+0.07}_{-0.07} \log \left( \frac{\Sigma_{\text{SFR},40}}{M_\odot \text{ kpc}^{-2} \text{ yr}^{-1}} \right) - 0.61^{+0.11}_{-0.12} \pm 0.19 \quad \text{Figure 8(h)}. \tag{63}
\]
The instantaneous SN-origin metal loading factor in cool outflows is nearly identical to that in hot outflows, slightly lower near the disk and higher farther away.

The metal enrichment factor $\zeta$ is nearly flat at low $\Sigma_{\text{SFR}}$, $\zeta \approx 1$ and 1.5 for cool and hot outflows, respectively. $\zeta$ begins to increase with $\Sigma_{\text{SFR}}$ above $\Sigma_{\text{SFR}} > 0.1 M_{\odot} \text{kpc}^{-2} \text{yr}^{-1}$, reaching $\zeta \approx 1.1$ and 2 for cool and hot outflows, respectively.

There is a very tight, positive correlation between energy and SN-origin metal fluxes (and hence loading factors) in the hot outflow. A similar but looser correlation also exists for the total outflow. Taking all outflow time series into account, we find that the correlation is slightly superlinear with a log–log slope of 1.15 at $|z| = 1 \text{kpc}$ (Equation (39)). This means that the energy loading in outflows is more efficient (radiative cooling is reduced) when more genuine SN material is loaded ($=$more successful breakout). This can be translated into a correlation between the Bernoulli velocity (or specific energy) and the SN-origin mass fraction in the outflow as in Equation (40).

6. Comparison with other simulations.—Our results are overall consistent with previous local simulations as long as their adopted $\Sigma_{\text{SFR}}$ are consistent with our predicted self-consistent values at a given $\Sigma_{\text{gas}}$ and vertical gravity. However, mass and energy partitions between phases may still be quite sensitive to the adopted SN distribution and its mutual correlation with gas distribution.

Finally, we close this paper by putting the present work in the context of the general goal of the SMAUG project: the development of physical subgrid models for galaxy formation models. Currently, in large-box cosmological simulations and SAMs, galactic winds are often implemented via scaling models. Currently, in large-box cosmological simulations and the context of the general goal of the SMAUG project: the relations quantified here provide represent unresolved galaxies of different mass at varying redshift.

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Software: Athena (Stone et al. 2008; Stone & Gardiner 2009), astropy (Astropy Collaboration et al. 2013, 2018), scipy (Virtanen et al. 2020), numpy (van der Walt et al. 2011), IPython (Perez & Granger 2007), matplotlib (Hunter 2007), linmix (Kelly 2007), xarray (Hoyer & Hamman 2017), pandas (Wes McKinney 2010), CMasher (van der Velden 2020), corner (Foreman-Mackey 2016), adstex (https://github.com/yymao/adstex).
Appendix A
Convergence with Resolution and Box Size

The numerical convergence of the TIGRESS framework has been extensively demonstrated and discussed in KO17 (for general ISM properties, SFRs, and outflow fluxes) and in KO18 (for multiphase characteristics of outflows). For the solar-neighbourhood model (R8 in Table 1), we showed that marginal convergence is achieved at 16 pc and more robust convergence at 8 pc. Due to the generally shorter dynamical timescales and length scales, we anticipate more stringent convergence conditions in higher-density environments. The simulation parameters shown in Table 1 indeed adopt finer spatial resolution for these models.

To test resolution convergence, here we present results from a model suite with two times poorer spatial resolution than the standard model suite. Note that given the stochastic nature of each simulation’s evolution, only statistical comparisons are possible between different resolutions. Figure A1 plots the mass loading factor of cool outflows and the energy loading factor of hot outflows for selected galactic properties (see Figures 9 and 10). The lower-resolution models are in good agreement with higher-resolution models, falling on the reported scaling relations within 1σ uncertainty levels. To test box size convergence, we rerun model R2 at lower resolution \( \Delta x = 4 \) pc and varying horizontal domain sizes from smaller \( L_x = L_y = 256 \) pc to standard 512 pc to larger 1024 pc, while our standard choice is \( \Delta x = 2 \) pc and \( L_x = L_y = 512 \) pc. We use model R2 because this model is expected to show the largest box size dependence owing to its shortest gravitational timescale comparable to star cluster evolution timescale. Figure A2 compares time evolution (left) and mean/standard deviation (right) over \( 0.5 < t/t_{\text{orb}} < 1.5 \) for a few selected quantities, \( \Sigma_{\text{SFR}} \) in panels (a) and (b), mass flux and loading of cool outflows in panels (c) and (d), respectively, and energy flux and loading of hot outflows in panels (e) and (f), respectively. We confirm that the lower-resolution model is in good agreement with the standard model as already demonstrated in Figure A1.

There are general increasing trends with box size in both mass and energy fluxes and loading factors. Our choice of box size is smaller than or comparable to the Toomre length scale,

\[
\lambda_T = \frac{4\pi^2G\Sigma_{\text{gas}}}{\kappa^2} = 850 \text{ pc} \left( \frac{\Sigma_{\text{gas}}}{100 M_\odot \text{ pc}^{-2}} \right) \left( \frac{\Omega}{100 \text{ km s}^{-1} \text{ kpc}^{-1}} \right)^{-2},
\]

above which axisymmetric gravitational instability is suppressed by epicyclic motions. This means that if the large-scale coherent structure is not destroyed by feedback within the gravitational timescale, the entire gas disk would collapse globally. For R2, \( t_f \lesssim t_{\text{osc}}, t_{\text{evol}} \), we anticipate large-scale gravitational collapse from the initial conditions. In this case, star formation is more clustered with a larger box, resulting in stronger feedback and higher loading factors, especially, for the energy loading factor (Figure A2(f)). Such strong bursts may indeed exist in galactic centers. As the validity of the local approximation is in question as \( L \) gets closer to \( R_0 \), however, we limit our model to a moderate box size, but still large enough to capture spatial correlation of SNe to some extent. Global modeling is clearly necessary in this regime.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figs/figA1.png}
\caption{Resolution convergence for scaling relations of loading factors. Top: scaling relations of cool mass loading factor. Bottom: scaling relations of hot energy loading factor. The mass and energy fluxes are both measured at \( [z] = H \). From left to right, the x-axes denote SFR surface density with \( \Sigma_{\text{SFR}} = 40 \text{ Myr}^{-1} \) (a)/(e), midplane total pressure \( (b)/(f)) \), total gas weight \( (c)/(g)) \), and gas depletion time \( (d)/(h)) \). Simulation results from standard and low-resolution model suites are presented as lighter and darker symbols, respectively (see legends distributed over panels). The best-fit lines for standard and low-resolution models are shown as black and magenta solid lines, respectively.}
\end{figure}
Appendix B

Instantaneous Loading Factors with Delayed Normalization

In our simulation suite (see, e.g., Figure 4), we observe more than an order of magnitude temporal fluctuations, and generally a delay between a peak in the SN rate and the enhancement in the outflow flux. As we discussed in Section 4.3, the complicated quantitative behavior makes it difficult to define instantaneous loading factors in realistic simulations where both feedback injection rates and outflow rates are self-consistently modulated (e.g., Muratov et al. 2015). For example, Figure B1 plots the normalized mass outflow rate and SN rate for all models. Overall, there is stronger temporal fluctuation in outflow rates than SN rates. Often, a moderate level of continuous SN explosions does not create corresponding outflows (e.g., $t/t_{\text{turb}} = 0.8$–1 for LGR8), mainly due to strong inflows of material ejected by previous outflows. For this reason, attempting a one-to-one mapping of the peaks of outflow rate and SN rate (or SFR) with a constant time delay generally fails in our simulations. It is worse at higher SFRs and not particularly better for different physical quantities (momentum, energy, and metal) and phases.

Nevertheless, we have tested defining a delay time in two ways, in order to investigate the uncertainty in calculation of loading factors. (1) We calculate the Pearson correlation coefficient between the SN rate and outflow rates averaged over $0.5 < t/t_{\text{turb}} < 1.5$. Right: mean and standard deviation over $0.5 < t/t_{\text{turb}} < 1.5$ for (b) $\Sigma_{\text{SFR,40}}$, (d) $\eta_{M,\text{cool}}$, and (f) $\eta_{E,\text{hot}}$.

Figure A2. Box size convergence test for R2. The “standard” model adopts a spatial resolution $\Delta x = 2$ pc and horizontal domain size $L_x = L_y = 512$ pc, while the other models adopt $\Delta x = 4$ pc with varying $L_x = L_y = L$ pc shown in the model name. Left: time evolution of (a) $\Sigma_{\text{SFR,40}}$, (c) $M_{\text{cool}}$, and (e) $M_{\text{cold}}$ over $0.5 < t/t_{\text{turb}} < 1.5$. Right: mean and standard deviation over $0.5 < t/t_{\text{turb}} < 1.5$ for (b) $\Sigma_{\text{SFR,40}}$, (d) $\eta_{M,\text{cool}}$, and (f) $\eta_{E,\text{hot}}$.

Table B1 lists the delay times and loading factors obtained by two methods, along with the loading factors without time delay. In Column (3). The delay times found in this way are longer in models with longer $t_{\text{osc}}$. The derived loading factors are consistent within the intrinsic uncertainty arising from the temporal fluctuations and mismatch between outflow and SN rates. The mass loading factor estimated by model fitting gives generally smaller values, but not very different from other estimates.
Figure B1. Comparison between outflow rate and SN rate. All quantities are normalized by their own mean over the time range shown in the plot. Mass outflow rates averaged over $|z| = H - 2H$ are compared with original and delayed SN rates. The time delay maximizing the Pearson correlation coefficient (Column (3) in Table B1) is applied.
Appendix C

Scaling Relations with Input Parameters

In the main portion of the paper, we provided scaling relations for loading factors with respect to the self-regulated ISM properties such as $S_{\text{FR}}$ and $P_{\text{mid}}$ (see Section 5). Here we additionally present scaling relations with respect to the input model parameters in Table 1. Figure C1 shows the mass loading factor of the cool outflow (top row) and the energy loading factor of the hot outflow (bottom row) at $|z| = H$ as a function of initial gas surface density ($\Sigma_{\text{gas,0}}$), stellar+dark matter midplane volume density ($\rho_{\text{sd}} \equiv \Sigma_*/(2z_*) + \rho_{\text{dm}}$), and angular velocity of galactic rotation ($\Omega$). Based on the intrinsic scatter ($\sigma_{\text{int}}$ in each panel of Figure C1), we find that the energy loading $\eta_{E,\text{hot}}$ correlates better with “gravity parameter” $\rho_{\text{sd}}$, while mass loading $\eta_{M,\text{cool}}$ correlates better with “gas parameter” $\Sigma_{\text{gas,0}}$. Both loading factors show good correlation with $\Omega$. Overall, $\eta_{M,\text{cool}}$ better correlates with the self-regulated ISM properties shown in Figure 9 than with the input model parameters shown in Figure C1. We note that the input parameters are not chosen to be fully independent of each other: roughly, $\Sigma_{\text{gas,0}} \propto \rho_{\text{sd}}$ with two different normalizations for the R and LGR series, and $\rho_{\text{sd}} \propto \Omega^2$.

Table B1

| Model | $(q,\text{ph})$ | $\eta_{q,\text{ph}}$ | $d_{\text{corr}}^q$ (Myr) | $\eta_{q,\text{corr}}$ | $d_{\text{model}}^q$ (Myr) | $\eta_{q,\text{model}}$ |
|-------|-----------------|------------------|-----------------|------------------|-----------------|------------------|
| R2    | $(E,\text{hot})$ | 0.15             | 1.4             | 0.14             | 2.0             | 0.14             |
| R4    | $(E,\text{hot})$ | 0.13             | 2.9             | 0.12             | 2.9             | 0.11             |
| R8    | $(E,\text{hot})$ | 0.053            | 16              | 0.053            | 16              | 0.071            |
| R8    | $(M,\text{cool})$ | 4.5              | 12              | 4.5              | 12              | 4.0              |
| R16   | $(E,\text{hot})$ | 0.023            | 21              | 0.023            | 21              | 0.025            |
| R16   | $(M,\text{cool})$ | 31               | 26              | 32              | 35              | 25               |
| LGR2  | $(E,\text{hot})$ | 0.074            | 4.9             | 0.07            | 4.9             | 0.079            |
| LGR2  | $(M,\text{cool})$ | 1.2              | 3.4             | 1.1             | 2.9             | 0.89             |
| LGR4  | $(E,\text{hot})$ | 0.057            | 3.9             | 0.058           | 2.9             | 0.063            |
| LGR4  | $(M,\text{cool})$ | 3.4              | 2.0             | 3.4             | 2.0             | 3.2              |
| LGR8  | $(E,\text{hot})$ | 0.053            | 11              | 0.055           | 11              | 0.089            |
| LGR8  | $(M,\text{cool})$ | 7.3              | 11              | 7.6             | 11              | 8.9              |

Note. Column (2): combination of the outflow quantity “$q$” and phase “$\text{ph}$” used to measure the delay time and loading factor. Column (3): loading factors reported in Section 5 without any time delay. Columns (4) and (5): delay time and loading factor maximizing the Pearson correlation coefficient between SN rate and outflow rate. Columns (6) and (7): delay time and loading factor minimizing the difference between the delayed and measured model fluxes.
Figure C1. Scaling relations of cool mass loading and hot energy loading factors with simulation input parameters. The mass and energy fluxes are measured at $|z| = z_{\text{H}}$. Figures at different heights are available at https://changgoo.github.io/tigress-wind-figureset/figureset.html. The simulation results and fitting results are presented as in Figure 8.

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