Dynamics of hot galactic winds launched from non-uniform starburst cores

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ABSTRACT

The galactic wind model of Chevalier & Clegg (1985) (CC85) assumes uniform energy and mass-injection within the starburst galaxy nucleus. However, the structure of nuclear star clusters, bulges, and star-forming knots are inherently non-uniform. We generalize to cases with non-uniform energy/mass injection that scale as \( r^{-\Delta} \) within the starburst volume \( R \), providing solutions for \( \Delta = 0, 1/2, 1, 3/2, \) and 2. In marked contrast with the CC85 model (\( \Delta = 0 \)), which predicts zero velocity at the center, for a singular isothermal sphere profile (\( \Delta = 2 \)), we find that the flow maintains a constant Mach number of \( M = \sqrt[3]{\frac{3}{5}} \approx 0.77 \) throughout the volume. The fast interior flow can be as \( v_{r<R} = (\frac{E_T}{3M_f})^{1/2} \approx 0.41 \times v_{\infty} \), where \( v_{\infty} \) is the asymptotic velocity, and \( E_T \) and \( M_f \) are the total energy and mass injection rates. For \( v_{\infty} \approx 2000 \) km s\(^{-1} \), \( v_{r<R} \approx 820 \) km s\(^{-1} \) throughout the wind-driving region. The temperature and density profiles of the non-uniform models may be important for interpreting spatially-resolved maps of starburst nuclei. We compute velocity resolved spectra to contrast the CC85 and \( \Delta = 2 \) models. Next generation X-ray space telescopes such as XRISM may assess these kinematic predictions.

Key words: galactic winds – starburst galaxies – hydrodynamics – X-rays:general

1 INTRODUCTION

Galactic winds are important to the process of galaxy formation and evolution (see Veilleux et al. 2005; Zhang et al. 2018; Veilleux et al. 2020). They are commonly found in rapidly star-forming galaxies at both low and high redshift (Martin 2005; Rubin et al. 2014), to modulate star formation, shape the stellar mass and metallicity relations (Peeples & Shankar 2011; Ma et al. 2016), and advect metals into the circumgalactic and intergalactic medium (Borthakur et al. 2013; Werk et al. 2016).

Galactic outflows are observed to be multi-phase. The hot, \( T \geq 10^7 \) K, phase is observed in X-rays and is often compared to the CC85 wind model (e.g., Strickland & Heckman 2009; Lopez et al. 2020, 2022). The CC85 model assumes uniform energy and mass-injection within a sphere of radius \( R \sim 0.1 \) – 0.5 kpc, which drives a flow that has the characteristic solution of transitioning from sub to supersonic at \( R \). Outside of the sphere, the flow undergoes adiabatic expansion (i.e., \( T \propto r^{-4/3}, \rho \propto r^{-2}, \) and \( v \propto r^0 \)).

There have been many modifications to CC85. These semi-analytic studies typically relax the assumption of an adiabatic wind by including additional physics such as radiative cooling, gravity, radiation pressure, non-equilibrium ionization, non-spherical flow geometries, and/or mass-loading of swept up material (see Wang 1995; Suchkov et al. 1996; Silich et al. 2004; Thompson et al. 2016; Boustead et al. 2016; Yu et al. 2020; Nguyen & Thompson 2021; Fielding & Bryan 2022; Sarkar et al. 2022). Other studies have numerically considered uniform wind-driving cylinders (Strickland et al. 2000) and rings (Nguyen & Thompson 2022), and non-uniform injection within cold galactic disks (Tanner et al. 2016; Schneider et al. 2020).

Star formation is inherently non-uniform. Embedded stellar clusters display either multi-peaked surface density distributions or highly concentrated surface density distributions (Lada & Lada 2003). Nuclear star clusters and bulges are observed to be compact and non-uniform (Böker et al. 2002). Consequently, a self-consistent wind model needs to consider non-uniform sources within the wind-driving region (WDR).

In this work, in contrast with uniform injection, we consider volumetric energy and mass injection, \( q \) [ergs cm\(^{-3} \)] and \( Q \) [g cm\(^{-3} \)] respectively, that scales as \( q \propto Q \propto r^{-\Delta} \) within the WDR. Zhang et al. (2014) present the solutions for arbitrary \( \Delta \) models but do not present a study on the bulk gas dynamics and thermodynamics of these models. Silich et al. (2011) presents wind models with non-uniform mass and energy injection modeled with an exponential function as \( q \propto Q \propto \exp(1-R/R_f) \). Both Palouš et al. (2013) and Boustead et al. (2016) consider a Schuster distribution of sources that scale as \( q \propto Q \propto (1-r^2/R_f^2)^\zeta \) (with the latter reference taking \( \zeta = 0 \)). In these previous works the sonic point shifts away from \( R \), as \( q \) and \( Q \) are taken to be non-uniform.
Here we calculate the structure of \( r^{-\Delta} \) models for \( r < R \). Similar to CC85, we assume that the supernovae efficiently thermalize their energy and drive a wind. We extend Zhang et al. (2014) by exploring how the kinematic and thermodynamic properties of the flow change over different injection slopes \( \Delta \). The CC85 model (\( \Delta = 0 \)) predicts flat temperature, densities, and pressure within the WDR and zero velocity at the center which linearly accelerates to become supersonic at the starburst ridge. We find the non-uniform models produce flows that are denser and faster than the CC85 flows within \( R \). Notably, for a model representative of a galactic density profile with a constant rotation curve, an isothermal sphere \( (\rho_{\text{sources}} \propto r^{-2}) \), the outflow maintains \( M = \sqrt{3/5} \approx 0.77 \) flow throughout the WDR. We verify these results using 3D hydrodynamic simulations with the Cholla code for \( \Delta = 1/2, 1, 3/2, \) and 2 models. We then focus on the observational characteristics of these non-uniform injection wind models, finding that the fast subsonic winds \( (v_r < R \approx 0.4 v_{\infty}, \text{see Eq. 15}) \), leads to horn-like features in resolved line profiles (Fig. 3) which may be observed by XRISM.

In §2, we write down the hydrodynamic equations, derive the self-similar analytic Mach number, physical, dimensionless solutions, and take central limits of these solutions. In §3, we run 3D hydrodynamic simulations, confirm the derived analytics, and construct X-ray surface brightness, brightness vs. height profiles, and velocity resolved line profiles. In §5, we provide a synthesis of this work, discuss how the models predict outflow velocities that can be resolved by XRISM, how the different \( T \) and \( n \) profiles may be important in interpreting spatially-resolved maps for the interior of starburst superwinds, and consider future research directions.

2 HYDRODYNAMIC EQUATIONS

In the absence of gravity and radiative cooling, the hydrodynamic equations for a steady-state spherically expanding flow are (see Chevalier & Clegg 1985):

\[
\begin{align*}
\frac{1}{r^2} \frac{d}{dr} (r^2 v_r^2) &= q, \\
\frac{dv}{dr} &= \frac{1}{\rho} \frac{dP}{dr} - \frac{q v}{\rho}, \\
\frac{1}{r^2} \frac{d}{dr} \left( r^2 v_r^2 \left( \frac{1}{2} v_r^2 + \frac{\gamma}{\gamma-1} \frac{P}{\rho} \right) \right) &= Q,
\end{align*}
\]

where the volumetric energy and mass injection rates are

\[
q = \begin{cases} 
q_0 (R/r)^\Delta, & (r \leq R) \\
0, & (r > R)
\end{cases}, \quad \text{and} \quad Q = \begin{cases} 
Q_0 (R/r)^\Delta, & (r \leq R) \\
0, & (r > R)
\end{cases},
\]

respectively. Equations 1, 2, and 3 can be re-written as a single equation, the derivative of the Mach number, as

\[
\frac{dM}{dr} = \frac{1}{(\gamma P/\rho)^{1/2}} \frac{d}{dr} \left( \frac{v}{(\gamma P/\rho)^{1/2}} \frac{1}{2 \rho} \frac{dP}{dr} - \frac{v}{(\gamma P/\rho)^{1/2}} \frac{1}{2P} \frac{dP}{dr} \right).
\]

We then impose the boundary condition (Chevalier & Clegg 1985; Wang 1995):

\[
M(r = R) = 1.
\]

For \( 1 \neq \Delta < 3 \), the solution for the Mach number as a function of radius within the WDR is

\[
M(1/(1-\Delta) \left[ \frac{2 + M^2 (y - 1)}{y + 1} \right] ^{(\Delta - 1) y + 1} \right) \frac{1}{M^2 (y - 1) + 1} = \frac{r - R}{R} \quad (r \leq R)
\]

For \( \Delta = 1 \), the solution is

\[
M(1/(1-\Delta) \left[ \frac{2 + M^2 (y - 1)}{y + 1} \right] ^{(\Delta - 1) y + 1} \right) \frac{1}{M^2 (y - 1) + 1} = \frac{r - R}{R} \quad (r \leq R)
\]

These solutions agree with those also derived by Zhang et al. (2017). Taking \( \Delta = 0 \), we arrive at the CC85 solution. The adiabatic, spherically expanding, exterior \( (r > R) \) solutions for the Mach number are identical to the CC85 solutions for all \( \Delta \) models.

Physical solutions require the definition of the total energy and mass injection rates, \( \dot{E}_T \) and \( \dot{M}_T \), within the WDR. We use

\[
\dot{M}_T = \beta \dot{M}_{\text{SFR}} \quad \text{and} \quad \dot{E}_T = \alpha \cdot \tau \times 10^{44} \dot{M}_{\text{SFR}} \quad [\text{ergs s}^{-1}]
\]

where \( \dot{M}_{\text{SFR}} = \dot{M}_{\text{SFR}} / M_{\odot} \text{yr}^{-1} \) is the dimensionless star-formation rate, and we have assumed that there is one supernova per 100 \( M_{\odot} \) of star formation and that each supernova releases \( 10^{51} \) ergs of energy. To make a direct comparison with the uniform case \( (\Delta = 0) \), we normalize the rates to that of the uniform CC85 model. The normalization requirement for energy and mass loading, with rates \( q' \propto \dot{Q}' \propto r^{-\Delta} \), are

\[
\dot{q}' = \frac{3 - \Delta}{3} \frac{\dot{M}_T}{(4/3) \pi R^3} \quad \text{and} \quad \dot{Q}' = \frac{3 - \Delta}{3} \frac{\dot{E}_T}{(4/3) \pi R^3}, \quad (\Delta < 3)
\]

From Equations 1 and 3, the sound speed and velocity are

\[
\frac{c_s^2}{\dot{M}_T} = \left( \frac{\dot{E}_T}{\dot{M}_T} \left[ \frac{(y - 1)M^2 + 2}{2(y - 1)} \right] \right)^{1/2} \quad \text{and} \quad v = \dot{M}_T \left[ \frac{\dot{E}_T}{\dot{M}_T} \left[ \frac{(y - 1)M^2 + 2}{2(y - 1)} \right] \right]^{1/2}.
\]

The density is obtained from the continuity equation as

\[
\rho = \frac{\dot{M}_T}{4 \pi v R^{3-\Delta}}, \quad (r \leq R) \quad \text{and} \quad \rho = \frac{\dot{M}_T}{4 \pi v R^2}, \quad (r > R).
\]

The remaining quantity, the pressure, is solved from the sound speed as \( P = \rho c_s^2 / \gamma \), where we take \( \gamma = 5/3 \) throughout the paper.

In Figure 1, we plot the dimensionless Mach number, density, and temperature. Relative to uniform injection \( (\Delta = 0, \text{red line}) \), we find an isothermal sphere model \( (\Delta = 2, \text{blue line}) \) produces a higher Mach number and denser outflow within the interior of the starburst. In Table 1 we present analytic central limit solutions for \( \Delta = 0 \) and 2 models. We find the Mach number for an isothermal sphere \( (\Delta = 2) \) is constant:

\[
M = \sqrt{3/5} \approx 0.77, \quad (\Delta = 2, r < R).
\]

This starkly contrasts the Mach number for a CC85 (uniform) model, which linearly grows as \( M = 2^{-5/14} 3^{-11/14} r_s \approx 0.33 r_s \) from the origin, where \( r_s = r / R \). Both the pressure and density scale as \( r^{-1} \) such that the velocity profile is also constant within the WDR:

\[
v = \left( \dot{E}_T / 3 \dot{M}_T \right)^{1/2}, \quad (\Delta = 2, r < R).
\]

This can be written in terms of the asymptotic wind-velocity, \( v_{\infty} = (2 \dot{E}_T / 3 \dot{M}_T)^{1/2} \), as

\[
v = v_{\infty} / \sqrt{6} = 0.4 v_{\infty}, \quad (\Delta = 2, r < R)
\]

We see that the interior flow for a \( \Delta = 2 \) model is approximately half
Analytic Central Solutions for Figure 1.

See Section 2 for central limit of the $\Delta = 1$ model. The physical variables are written in cgs units and assume $\mu = 0.6$.

| Dimensionless Variable | Uniform Sphere | Isothermal Sphere |
|------------------------|----------------|-------------------|
| $M$                    | $r_*/(25^{1/4}4^{1/4})$ | $\sqrt{575}$ |
| $v/M_T^{-1/2}E_T^{1/2}$ | $21^{1/3}3^{-9/7}r_*$ | $1/\sqrt{3}$ |
| $T/M_T^{-1}E_T\mu p_k^{-1}$ | $2/5$ | $1/3$ |
| $\rho/M_T^{-3/2}E_T^{-1/2}R^{-2}$ | $3^{9/7}/(4\pi 2^{7/7})$ | $\sqrt{3}/(4\pi r_*)$ |
| $P/M_T^{-1/2}E_T^{1/2}R^{-2}$ | $3^{9/7}/(5\pi 2^{9/7})$ | $1/(4\pi \sqrt{3}r_*)$ |

Physical Variable (Solar metallicity, cgs units)

| Uniform Sphere | Isothermal Sphere |
|----------------|-------------------|
| $r_*/(25^{1/4}4^{1/4})$ | $\sqrt{575}$ |
| $v/M_T^{-1/2}E_T^{1/2}$ | $21^{1/3}3^{-9/7}r_*$ | $1/\sqrt{3}$ |
| $T/M_T^{-1}E_T\mu p_k^{-1}$ | $2/5$ | $1/3$ |
| $\rho/M_T^{-3/2}E_T^{-1/2}R^{-2}$ | $3^{9/7}/(4\pi 2^{7/7})$ | $\sqrt{3}/(4\pi r_*)$ |
| $P/M_T^{-1/2}E_T^{1/2}R^{-2}$ | $3^{9/7}/(5\pi 2^{9/7})$ | $1/(4\pi \sqrt{3}r_*)$ |

as fast as the energy-conserving supersonic terminal velocity. The difference in kinematics may be observable in velocity-resolved line profiles (see Sec. 4).

For $\Delta = 1$, in the limit that $r_*/r < 1$, the Mach number is given by $M = [-20/3 \times \ln((4/3)^{1/5}r_*)]^{1/2}$. The remaining quantities may be calculated by combining this with Equations 11 and 12.

Equations 7 and 8 are solutions to an implicit equation. To use the solution, one is required to define an inner and outer radius. As shown in Table 1, for an isothermal sphere model, there is a strong dependence on the inner radius, as the density and pressure diverge towards infinity (i.e., $\rho_*/p_* \propto r_*^{-1}$ for $\Delta = 2$ and $r_*/r < 1$). We define the inner radius for the $\Delta = 2$ model as $r_{core,\infty}$.

2.1 Inference of the volumetric energy and mass injection rates within the wind-driving region

A critical inference from X-ray observations of starburst nuclear centers are the energy thermalization and mass-loading efficiencies (i.e., $\alpha$ and $\beta$, see Eqs. 9). Using measurements of the central temperature and density, we infer $\alpha$ and $\beta$ (Strickland & Heckman 2009) using the solutions from Table 1, for both $\Delta = 0$ and $\Delta = 2$. These efficiencies are derived to be

$$\alpha(n,T) = \begin{cases} 0.105 n_0.1 T_7^{3/2} R_9^{5/2} M_5^{-1}, & (\Delta = 0) \\ 0.297 n_0.1 T_7^{3/2} R_9^{5/2} M_5^{-1} r_{core,\infty}, & (\Delta = 2) \end{cases}$$

and

$$\beta(n,T) = \begin{cases} 0.150 n_0.1 T_7^{3/2} R_9^{5/2} M_5^{-1}, & (\Delta = 0) \\ 0.353 n_0.1 T_7^{3/2} R_9^{5/2} M_5^{-1} r_{core,\infty}, & (\Delta = 2) \end{cases}$$

where $n_0.1 = n/0.1 \text{ cm}^{-3}$, $T_7 = T/10^7 \text{ K}$, $R_9 = R/0.5 \text{ kpc}$, and $M_5 = M_{\text{SFR,}\infty}/10$. From Equations 16 and 17, it is apparent that inferred efficiencies $\alpha$ and $\beta$ from the $\Delta = 2$ model have a dependence on the defined inner radius $r_{core,\infty}$, whereas for $\Delta = 0$, there is not.

3 3D HYDRODYNAMIC SIMULATIONS

We test our solutions using the Cho11a (Schneider & Robertson 2015) code to simulate the starburst nuclei. Each simulation is carried out in a cube with uniform grid cells. The box has dimensions 1 kpc with 256$^3$ cells, giving a cell resolution of $\Delta_\perp = 3.9$ pc.

Within a radius of $R$, we deposit energy and mass at a rate $E_T$ and $M_T$ for different power-law injection slopes $\Delta = 0, 1/2, 1, 3/2, 2$, while the normalization for each $\Delta$ model is given by Equation 10. For all simulations, we take the M82-like fiducial wind parameters (Strickland & Heckman 2009) of $\alpha = 1$, $\beta = 0.3$, $R = 0.3 \text{ kpc}$, and $M_{\text{SFR,}\infty} = 10$. The value of the core radius $r_{core,\infty}$ is effectively set by the resolution. In order to make a direct comparison between the analytic solutions and numerical simulations, we do not include any additional physics, such as radiative cooling or gravity. For these wind model parameters, most of the flow is non-radiative and can escape a typical potential (Chevalier & Clegg 1985; Thompson et al. 2016; Lochhaas et al. 2021). All wind models reach a steady state, showing that the solutions are stable.
In Figure 2, we show 1D radial profiles (+z skewers) of the Mach number, number density, temperature, and velocity profiles for both the analytic solutions (colored solid lines) and the Cholla simulations (solid black markers) after a time-steady solution has been established. The analytic solutions match the simulation results for every physical quantity. This implies that the imposed boundary condition of $M = 1$ at $r = 1$, which was used in the analytic derivation, is indeed valid over the range of $\Delta$ values presented. Compared to the uniform sphere CC85 model ($\Delta = 0$), the isothermal sphere model ($\Delta = 2$) maintains a much higher, constant, radial velocity $v / 740 \text{ km s}^{-1}$ throughout most of the WDR.

4 OBSERVATIONAL SIGNATURES

4.1 Surface Brightness

We calculate the instantaneous X-ray surface brightness as $S_X(r, z) = \int_0^\infty d\nu \int_0^{L_S} n(x, y, z) \Delta(T(x, y, z), \nu, z)$, where $L_S$ is the length of the simulation domain, which includes the post-WDR supersonic wind. Using PyAtomDB (Foster & Heuer 2020), we evaluate the plasma emissivity over XRISM’s observing bandwidth ($0.3 \leq E_{\gamma} [\text{keV}] \leq 12$), and assume solar metallicity abundances (Anders & Grevesse 1989). In Figure 3 we show $S_X$ for $\Delta = 0, 1, 2$ Cholla models. In the left panel of Figure 4, we calculate the surface brightness profile by taking $\Sigma_X$ and integrating along the $\hat{z}$ direction, and then dividing by the area of each surface. The surfaces are taken to be $50 \text{ pc}^2$. The $\Delta = 2$ model leads to a strongly-peaked brightness profile, whereas the $\Delta = 0$ model produces a less-peaked profile. We note that for the $\Delta = 2$ model, the diverging density (see Tab. 1) implies a short cooling timescale. For these short cooling times, a cool non-X-ray emitting core may develop (Wünsch et al. 2008; Lochhaas et al. 2021). Radiative cores will be considered in a future work.

4.2 Velocity Resolved Line Profile

XRISM’s Resolve instrument is capable of resolving individual spectral lines and will trace gas motions through Doppler broadening and line shifts (XRISM Science Team 2020). A spectrum of the entire wind-driving region will yield insight into the hot gas kinematics, which remain thus far unprobed. We construct resolved velocity line profiles for the $\Delta = 0$ and $\Delta = 2$ wind models. To do so, we consider shells inside $r \leq R$. When projected along the line of sight, this leads to a top-hat distribution in $n(r)$ versus $v(r)$ space, with bounds defined by $\pm v_s(r)$. We then calculate the emissivity of O vii, Mg xi, and Si xii, and integrate over XRISM’s observing bandwidth. Next, we integrate over the volume $V = \Delta r_i 4\pi r_i^2 \Delta r_i$. The result is shown in the three right panels of Figure 4. The $\Delta = 2$ model produces brighter emission along higher velocities, whereas the $\Delta = 0$ model is brightest where the gas is stationary. This is a result of the constant high velocity flow (Eq. 15). For these injection parameters (see Sec. 3) the characteristic feature of the $\Delta = 2$ model is the sharp increase in the emissivity at $v / 750 \text{ km s}^{-1}$.

5 SUMMARY

In this work we study the dependence of injection slope, $\Delta$, for the kinematic and thermodynamic structure of the wind within the wind driving region. We derive analytic solutions, present their limits at small $r$ (see Table 1), and then confirm them with 3D Cholla simulations (see Fig. 2). Importantly, we find that for a distribution of sources that scale as $r^{-2}$ ($\Delta = 2$) the Mach number in the WDR is constant ($M = \sqrt{3/5}$) and is approximately half of the asymptotic wind velocity (see Eq. 15), faster than the uniform distribution...
Figure 4. Left panel: The X-ray surface brightness as a function of height for each $\Delta$ model with photons of energies $0.3 \leq E [\text{keV}] \leq 12$. The $\Delta = 2$ model produces a strongly peaked surface X-ray brightness profile, whereas the the CC85 X-ray surface brightness appears more broadened within the WDR. Right three panels: The velocity resolved line profile with emissivities integrated over energies $0.3 \leq E [\text{keV}] \leq 12$, for three He-like triplets: O vii, Mg xi, and Si xiii, respectively. The emissivities are calculated with PyAtomDB with energies corresponding to the bandpass of XRISM’s Resolve soft X-ray spectrometer instrument. We see that $\Delta = 2$ models lead to a sharp horn-like feature in the velocity distribution across all He-like triplets, with the discrepancy between $\Delta = 2$ and $\Delta = 0$ more apparent in the heavier triplets.

(Chevalier & Clegg 1985) or Schuster-like distributions (Palouš et al. 2013; Bustard et al. 2016). The inferred energy and mass-loading efficiencies, $\alpha$ and $\beta$, are affected by $\Delta$, with $\Delta = 2$ sensitive to the core radius of injection. The $\Delta = 2$ model produces strongly peaked X-ray brightness profiles (see Fig. 3). Figure 4 shows resolved line velocity profiles for relevant emission lines O vii, Mg xi, and Si xiii for the $\Delta = 0$ (CC85), 1, and 2 (isothermal sphere) models. These features may be observed by XRISM in the future. The $T$ and $n$ structure of the non-uniform models may be important in interpreting spatially-resolved maps for the interior of starburst superwinds.

Wünsch et al. (2008); Lochhaas et al. (2021) showed that in cases of high mass-loading, the WDR develops a cool inert core. To make a direct comparison to the analytics, the simulations did not include cooling. We expect a cool core at the origin, as $\rho \propto r^{-1/2}$ ($\Delta = 2$). This would affect the X-ray surface brightness profiles shown in Section 4. The condition for a cool inert core depends on the competing cooling and advection timescales and will be investigated in a future work.

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REFERENCES

Anders E., Grevesse N., 1989, gca, 53, 197
Böker T., Laine S., van der Marel R. P., Sarzi M., Rix H.-W., Ho L. C., Shields J. C., 2002, AJ, 123, 1389
Borthakur S., Heckman T., Strickland D., Wild V., Schiminovich D., 2013, ApJ, 768, 18
Bustard C., Zweibel E. G., D’Onghe E., 2016, ApJ, 819, 29
Chevalier R. A., Clegg A. W., 1985, Nature, 317, 44
Fielding D. B., Bryan G. L., 2022, ApJ, 924, 82
Foster A. R., Heuer K., 2020, Atoms, 8, 49
Lada C. J., Lada E. A., 2003, ARA&A, 41, 57
Lochhaas C., Thompson T. A., Schneider E. E., 2021, MNRAS.
Lopez L. A., Mathur S., Nguyen D. D., Thompson T. A., Olivier G. M., 2020, ApJ, 904, 152
Lopez L. A., Lopez L. A., Nguyen D. D., Thompson T. A., Mathur S., Bolatto A. D., Vulic N., Sardone A., 2022, arXiv e-prints, p. arXiv:2209.09260
Ma X., Hopkins P. F., Faucher-Giguère C.-A., Zolman N., Muratov A. L., Kereš D., Quataert E., 2016, MNRAS, 456, 2140
Martin C. L., 2005, ApJ, 621, 227
Nguyen D. D., Thompson T. A., 2021, MNRAS, 508, 5310
Nguyen D. D., Thompson T. A., 2022, ApJ, 935, L24
Palouš J., Wünsch R., Martínez-González S., Hueyotl-Zahuantitla F., Silich S., Tenorio-Tagle G., 2013, ApJ, 772, 128
Peebles M. S., Shankar F., 2011, MNRAS, 417, 2962
Rubin K. H. R., Prochaska J. X., Koo D. C., Phillips A. C., Martin C. L., Winstrom L. O., 2014, ApJ, 794, 156
Sarkar K. C., Sternberg A., Gnat O., 2022, arXiv e-prints, p. arXiv:2203.15814
Schneider E. E., Robertson B. E., 2015, ApJS, 217, 24
Schneider E. E., Ostriker E. C., Robertson B. E., Thompson T. A., 2020, arXiv e-prints, p. arXiv:2002.10468
Silich S., Tenorio-Tagle G., Rodríguez-González A., 2004, ApJ, 610, 226
Silich S., Bisnovatyi-Kogan G., Tenorio-Tagle G., Rodríguez-González A., 2004, ApJ, 610, 226
Suchkov A. A., Berman V. G., Heckman T. M., Balsara D. S., 1996, ApJ, 463, 528
Tanner R., Cecil G., Heitsch F., 2016, ApJ, 821, 7
Thompson T. A., Quataert E., Zhang D., Weinberg D. H., 2016, MNRAS, 455, 1830
Strickland D. K., Heckman T. M., 2009, ApJ, 697, 2030
Strickland D. K., Heckman T. M., Weaver K. A., Dahlem M., 2000, AJ, 120, 2965
Suchkov A. A., Berman V. G., Heckman T. M., Balsara D. S., 1996, ApJ, 463, 528
Winstrom L. O., 2014, ApJ, 794, 156
Zhang D., Davis S. W., Jiang Y.-F., Stone J. M., 2018, ApJ, 854, 110
Zhang D., Thompson T. A., Murray N., Quataert E., 2014, ApJ, 784, 93
Zhang D., Thompson T. A., Quataert E., Murray N., 2017, MNRAS, 468, 4801
Zhang D., Davis S. W., Jiang Y.-F., Stone J. M., 2018, ApJ, 854, 110

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