Dusty Cloud Acceleration with Multiband Radiation

Xiaoshan Huang, Shane W. Davis, and Dong Zhang

1 Department of Astronomy, University of Virginia, Charlottesville, VA 22904, USA
2 Department of Astronomy, University of Michigan, Ann Arbor, MI 48109, USA

Received 2019 August 2; revised 2020 March 17; accepted 2020 March 17; published 2020 April 15

Abstract

We perform two-dimensional and three-dimensional simulations of cold, dense clouds, which are accelerated by radiation pressure on dust relative to a hot, diffuse background gas. We examine the relative effectiveness of acceleration by ultraviolet (UV) and infrared (IR) radiation fields, both independently and acting simultaneously on the same cloud. We study clouds that are optically thin to IR emission but with varying UV optical depths. Consistent with previous work, we find relatively efficient acceleration and long cloud survival times when the IR band flux dominates over the UV flux. However, when the UV flux is dominant or even a modest percentage (~5%–10%) of the IR irradiating flux, it can act to compress the cloud, first crushing it and then disrupting the outer layers. This drives mixing of the outer regions of the dusty gas with the hot diffuse background to the point where most dust is not likely to survive or stay coupled to the gas. Hence, the cold cloud is unable to survive for a long enough timescale to experience significant acceleration before disruption even though efficient IR cooling keeps the majority of the gas close to the radiative equilibrium temperature (T \text{IR} \lesssim 100 \text{K}). We discuss the implications for observed systems, concluding that radiation pressure driving is most effective when the light from star-forming regions is efficiently reprocessed into the IR.

Unified Astronomy Thesaurus concepts: Galactic winds (572); Radiative magnetohydrodynamics (209); Radiative transfer simulations (1967); Hydrodynamical simulations (767); Interstellar medium wind (848)

1. Introduction

Galactic outflows are observed ubiquitously in star-forming galaxies (Veilleux et al. 2005). In many cases, the observed outflow rates and velocities imply that the outflows must have an important impact on the evolution of galaxies. These outflows almost certainly play a role in regulating star formation (e.g., Benson et al. 2003; Bower et al. 2012; Bolatto et al. 2013; Puchwein & Springel 2013), but they also affect the chemical evolution of galaxies as well as the circumgalactic and intergalactic medium (Aguirre et al. 2001; Erb 2008; Finlator & Davé 2008; Oppenheimer et al. 2010).

Multiphase winds have been observed in many star-forming galaxies, including molecular gas (e.g., Veilleux et al. 2009; Cicone et al. 2014; Leroy et al. 2015; Oosterloo et al. 2017; Walter et al. 2017; Geach et al. 2018; Spilker et al. 2018; Krieger et al. 2019), neutral atomic gas (e.g., Heckman et al. 1990; Rupke et al. 2002; Martin 2005; Morganti et al. 2016; Schulz et al. 2018), and ionized gas (e.g., Martin 1998; Shopbell & Bland-Hawthorn 1998; Muñoz-Vegara et al. 2019; Tadhunter et al. 2019). Although not observed directly, constraints on hot gas outflows are provided by X-ray observations of star-forming galaxies (Strickland & Heckman 2007; Zhang et al. 2014).

It has generally been believed that these outflows are driven primarily by supernova feedback. In high-star-formation-rate galaxies, the overlapping supernova remnants merge to form giant bubbles of hot gas that break out of the cooler interstellar medium (ISM) gas. The cooler gas is then entrained in these hot outflows (Chevalier & Clegg 1985; Strickland & Heckman 2009). However, it is not entirely clear that the much cooler gas will survive long enough to reach the inferred velocities due to shredding and mixing with the hotter background flow (Begelman & Fabian 1990; Klein et al. 1994; Scannapieco & Brüggen 2015; Zhang et al. 2018). One possibility is that magnetic fields suppress the disruption due to purely hydrodynamical instabilities (McCourt et al. 2015; Banda-Barragán et al. 2016), but it is unclear if the magnetic-field strengths and geometries are present in such outflows. Another possibility is that the gas is destroyed but condenses out of the hot flow due to radiative cooling at larger distances from the galaxy (Thompson et al. 2016).

It is possible that a number of different acceleration mechanisms play a role in launching outflows, with different mechanisms possibly dominating in different environments or at different stages in the acceleration (Hopkins et al. 2012; Zhang 2018). In addition to entrainment, plausible mechanisms include radiation pressure of starlight on dust (e.g., Murray et al. 2005, 2011), cosmic-ray pressure (e.g., Ipavich 1975; Socrates et al. 2008; Ruszkowski et al. 2017; Wiener et al. 2019), and active galactic nuclei (Fabian 2012; Heckman & Best 2014).

In this work, we focus on the role played by the radiation pressure of starlight on dust. This mechanism has already been extensively studied with detailed radiation hydrodynamic numerical simulations at different scales and with varying assumptions and setups. An important question has been the role of Rayleigh–Taylor instabilities in limiting the effectiveness of acceleration (Krumholz & Thompson 2012, 2013; Skinner & Ostriker 2015; Raskutti et al. 2016). Despite the presences of such instabilities, it seems that some fraction of the radiation can be accelerated to large velocities (Davis et al. 2014; Tsang & Milosavljević 2015; Zhang & Davis 2017) and may allow cold gas to survive longer than entrainment allows (Zhang et al. 2018). With the exception of Raskutti et al. (2016), who studied the effect of radiation pressure in the local star-cluster environment, most of these studies focus on infrared (IR) radiation pressure. This is sensible for considering the role of radiation pressure on galactic scales since the vast majority of the light in the most extreme star-forming galaxies...
(luminous or ultraluminous IR galaxies, hereafter LIRGs and ULIRGs) is reprocessed into the IR. However, most of the radiation originates from starlight radiated in the ultraviolet (UV) band. It is possible that the UV band plays a greater role in the launching of gas close to the star clusters or in galaxies that lower dust obscuration. Indeed, previous work has shown that UV radiation pressure may be important in the vicinity of massive star formation; although, its role is limited by inhomogeneities in the gas arising from Rayleigh–Taylor instabilities (Rosen et al. 2016) and inhomogeneities in the radiation field itself from multiple stellar sources (Kim et al. 2018). Although we neglect the impact of photoionization, it may also be important for the destruction of molecular gas near starbursts (e.g., Kim et al. 2016).

Therefore, we are motivated to consider the relative role played by UV and IR opacities in accelerating clouds.

This paper is organized as follows. In Section 2, we describe our numerical simulation methods and problem setup. In Section 3, we report on the results of a variety of simulations with differing assumptions about parameters of interest such as IR to UV flux ratios, optical depths, as well as the sensitivity to assumptions in the numerical method and simulation setup. We discuss the primary implications of our results in Section 4 and summarize our conclusions in Section 5.

2. Method

2.1. Radiation Hydrodynamics Equations

We solve the equations of hydrodynamic and radiation transfer (RT) using the Athena++ (J. M. Stone et al. 2020, in preparation) code. The relevant equations are, respectively, the equations of conservation of mass, momentum, and energy, as follows:

\[
\frac{\partial\rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \\
\frac{\partial(\rho \mathbf{v})}{\partial t} + \nabla \cdot (\rho \mathbf{v} \mathbf{v} + \mathbf{P}) = -\mathbf{G}, \\
\frac{\partial E}{\partial t} + \nabla \cdot [(E + P)\mathbf{v}] = -cG_0, 
\]

(1)

Here, \(\rho\), \(\mathbf{v}\), \(E\), and \(P\) are the fluid density, velocity, total energy density, and pressure, respectively, and \(\mathbf{P}\) is the pressure tensor. The source terms \(\mathbf{G}\) and \(G_0\) represent the components of the radiation four-force, which are calculated by taking moments of the RT equation. The total energy density is

\[
E = \frac{P}{\gamma - 1} + \frac{1}{2}\rho v^2, 
\]

(2)

where \(\gamma\) is the adiabatic index and the terms represent the gas internal energy and kinetic energy, respectively. The radiation four-force is computed from the specific intensity \(I_\nu\), which is evolved according to the time-dependent RT equation:

\[
\frac{\partial I_\nu}{\partial t} + c \mathbf{n} \cdot \nabla I_\nu = S_\nu(n). 
\]

(3)

Here, \(S_\nu(n)\) is the radiation source term, \(n\) represents a unit vector parameterizing the direction, and \(c\) is the speed of light. The RT equation is solved using an explicit–implicit scheme in Eulerian frame, similar to the method described in Jiang et al. (2014). The main difference is that Jiang et al. (2014) evaluate radiation source terms in the Eulerian frame by expanding to the second order in \(v/c\). In the Athena++ implementation, the specific intensities \(I_\nu\) are transformed to the fluid comoving frame, where the opacities and emissivities are the simplest. The relevant source terms are evaluated and updated implicitly along with the radiation frame gas internal energy equation. The resulting source terms are integrated over frequency and angle, and then they are transformed back to the Eulerian frame.

In this work, we integrate Equation (3) over frequency assuming that the radiation field can be approximated with two radiation bands representing the IR and optical/UV contributions to the radiation field. The resulting RT equations solved are

\[
\frac{1}{c} \frac{\partial I_{\nu}}{\partial t} + \mathbf{n} \cdot \nabla I_{\nu} = -\Gamma(n)\kappa_{\nu\nu}\rho I_{\nu}, \\
\frac{1}{c} \frac{\partial I_{\lambda}}{\partial t} + \mathbf{n} \cdot \nabla I_{\lambda} = \Gamma(n)\kappa_\lambda \rho \frac{a_\lambda T_{\lambda}^4}{4\pi} - I_{\lambda}, 
\]

(4)

where \(a_\lambda\) is the radiation constant. The subscripts \(\nu\) and \(\lambda\) label the ultraviolet and infrared opacities and radiation fields, respectively. Note that we have assumed that there is no source of UV emission within the domain and that the UV radiation only originates from an external source via the boundary conditions. \(\Gamma(n)\) accounts for transformations between the comoving and Eulerian frames and corresponds to

\[
\Gamma(n) = \gamma_L \left(1 - \frac{v}{c}\right). 
\]

(5)

where \(\gamma_L\) is the Lorentz factor. The specific intensities \(I_{\lambda}\) and \(I_{\nu}\) are evaluated in the Eulerian frame, and the opacities \(\kappa_{\lambda}\) and \(\kappa_{\nu\nu}\) are evaluated in the comoving frame. However, we emphasize that the differences between the comoving and Eulerian frames are quite small in these simulations. The radiation energy and momentum terms are computed by integrating the appropriate moments of Equation (4) over angles.

\[
G = \int d\Omega I_{\nu} \Gamma(n) \left[\kappa_{\lambda} \rho \frac{a_\lambda T_{\lambda}^4}{4\pi} - \kappa_{\nu\nu} \rho (I_{\lambda} + I_{\nu})\right], \\
G_0 = \int d\Omega I_{\lambda} \Gamma(n) \left[\kappa_{\lambda} \rho \frac{a_\lambda T_{\lambda}^4}{4\pi} - \kappa_{\nu\nu} \rho (I_{\lambda} + I_{\nu})\right]. 
\]

(6)

(7)

In the limit of zero velocity, the corresponding momentum and energy source terms are the familiar expressions

\[
G \rightarrow -\frac{\kappa_{\nu\nu}\rho}{c} F_{\nu\nu} + \frac{\kappa_\lambda \rho}{c} F_{\lambda}, \\
G_0 \rightarrow \kappa_\lambda \rho \frac{a_\lambda T_{\lambda}^4 - E_\lambda}{c} - \kappa_{\nu\nu} \rho E_{\nu\nu}, 
\]

(8)

Here, \(F_{\lambda}\) and \(F_{\nu\nu}\) are the IR and UV radiation flux, respectively, and \(E_\lambda\) and \(E_{\nu\nu}\) are the IR and UV radiation energy density, respectively.

The UV dust opacity \(\kappa_{\nu\nu}\) depends on dust grain sizes and species as well as frequency, but we use a constant representative value. The IR opacity \(\kappa_\lambda\) is assumed to be a temperature-dependent Rosseland mean opacity using the approximation of Krumholz & Thompson (2012). To focus on the cold cloud dynamics, we ignore the scattering opacity,
setting scattering opacity to zero. The dust opacity is

\[
\kappa_d = 10^{-3/2} \left( \frac{T}{10^2 \text{K}} \right)^2 \text{s cm}^2 \text{g}^{-1},
\]

\[
\kappa_{uv} = 100 \text{s cm}^2 \text{g}^{-1}.
\]

This assumes a Milky Way–like dust-to-gas ratio, and \( \kappa_d \) is a reasonable approximation for \( T \lesssim 10^2 \) K and flattens at higher temperatures (Semenov et al. 2003). Hence, we assume a constant value of \( 10^{-3/2} \text{cm}^2 \text{g}^{-1} \) for \( T > 10^2 \) K. The quantity \( s \) represents a scaled ratio of the dust-to-gas fraction, normalized so that \( s = 1 \) corresponds to the initial dust-to-gas fraction (assumed to be uniform) in the cloud. The UV dust opacity will be dependent on both the grain size distribution and the wavelength of the optical-to-UV spectral energy density of the source driving the cloud acceleration. Our fiducial value is probably somewhat conservative in that dust opacity at UV wavelengths can be several times higher than that in the Milky Way ISM (e.g., Draine 2011). In our calculation, this may be offset by considering clouds with larger sizes or densities to make them more optically thick, although this choice does impact the ratio of IR to UV opacity, which may be an interesting parameter to explore in future work.

In order to track dust evolution, we initialize a cold cloud gas with \( s = 1 \) and cells in the hot background gas with \( s = 0 \). We then evolve \( s \) as a passive scalar via a continuity equation

\[
\frac{\partial s}{\partial t} + \mathbf{v} \cdot \nabla s = 0,
\]

which assumes that there is no source of dust other than the initial dust in the cloud. However, we adopt a simple prescription to account for the decoupling and destruction of the dust when it mixes with the hotter, less-dense background gas, setting the passive scalar to zero for cells above a fiducial temperature of \( 1500 \) K. This is approximately the temperature where most grain constituents are destroyed (Pollack et al. 1994), resulting in drops in the opacity (Semenov et al. 2003). Mixed gas at these temperatures also typically has densities low enough that it becomes uncertain whether the dust remains dynamically well-coupled with the gas, due to the increase in the mean free path of dust–gas collisions (Krumholz \\& Thompson 2013). A more sophisticated model of grain–gas interaction and grain destruction will be of interest in future studies, but our simple scheme serves its primary purpose, which is to decouple hotter and more diffuse gas from the radiation field. We choose \( 1500 \) K as a representative value, but we have checked that decreasing our decoupling temperature to \( 500 \) K or \( 1000 \) K has no significant impact on our inferred survival times.

Since the transport portion of the transfer equation is solved explicitly, the Courant–Friedrichs–Lewy (CFL) condition in the code is set by the speed of light, which is much larger than the flow velocity or sound speed. Hence, it is advantageous to adopt the reduced speed-of-light approximation, where \( c \) in Equation (4) is replaced by \( \tilde{c} = Rc \). Assuming \( R \ll 1 \) allows one to take time steps that are a factor of \( R^{-1} \) larger. As long as \( R \) is not chosen to be too low, the time-dependent term remains small, and the radiation flux remains close to quasi-steady on the flow timescale.

The conditions for the validity of the reduced speed-of-light approximation are described by Skinner \\& Ostriker (2013), which we outline in Section 3.2. The main constraint is the need to preserve the correct ordering of characteristic timescales. The light-crossing (i.e., radiation diffusion) time should always be smaller than the dynamical time. The radiation travels at reduced light speed \( \min(\tilde{c}, c/\tau_{\text{max}}) \), where \( \tau_{\text{max}} \) is the maximum optical depth in the system. For a system with characteristic length \( l_0 \), \( l_0/\min(\tilde{c}, c/\tau_{\text{max}}) \ll l_0/\tau_{\text{max}} \), where the dynamical timescale is \( l_0/v_{\text{max}} \), and \( v_{\text{max}} \) is the velocity that determines the dynamical timescale. For the modest flow velocities and low optical depths considered here, these constraints are easily obeyed for \( R = 0.01 \).

2.2. Simulation Setup

We initialize all simulations with a cold dense cloud in pressure equilibrium with a hotter, less-dense background gas. The cloud geometry is circular (2D) or spherical (3D), and it is initialized at rest in the center of the domain. A summary of the simulation parameters is provided in Table 1. We initialize the cloud to fiducial temperature \( T_0 = 50 \) K. We define a corresponding fiducial flux \( F_0 = c a r_0^2 \). We first model the UV radiation from a galaxy or star-forming region within the galaxy as a constant uniform flux \( F_{\text{uv}} = 1.4F_0 \approx 4.9 \times 10^{12} L_\odot \text{kpc}^{-2} \), and we inject the radiation flux from the bottom boundary. We used an angular grid described by Bruls et al. (1999). For 2D simulations, we used six angles per octant, yielding 84 angles in total. For 3D simulations, we used four angles per octant, yielding a total of 80 angles. The radiation flux in our simulation is scaled to the luminosity of ultraluminous infrared galaxies (ULIRGs). This choice of \( F_{\text{uv}} \) is about an order of magnitude smaller than typical IR radiation flux from ULIRGs (Zhang et al. 2018), which is higher than typically observed. We choose a high value for our fiducial flux to provide favorable conditions for acceleration, but we consider lower values in other calculations. In Section 4, we discuss the impact of varying the radiation flux on cloud dynamics.

The incoming radiation field injected at the lower \( x \) boundary is isotropic for incoming radiation on this boundary to model a distributed source of UV and IR emission, as might be expected from a large starbursting region of a galaxy. This is probably a good approximation for the IR radiation, but the UV radiation may be more directed if the radiation field is dominated by a few relatively distant star clusters. The (half) isotropic radiation field allows for optically thick clouds to be compressed in directions both parallel and perpendicular to the motion. A more parallel directed radiation field might be expected to compress the cloud primarily along the direction of motion (Proga et al. 2014) leading to more of a “pancake” structure in the initial evolution.

Given \( T_0 \), the fiducial speed \( v_0 \) is chosen to be the adiabatic sound speed \( c_s^2 = k_b T_0/(\mu m_\text{H}) \), where \( k_b \) is the Boltzmann constant, and we assume the mean molecular weight \( \mu = 1.0 \) for simplicity. Zhang et al. (2018) showed that varying \( \mu \) has a limited impact on cloud dynamics. For our fiducial run, we choose a fiducial initial density \( \rho_0 = 1.0 \times 10^{-19} \text{g cm}^{-3} \). We applied a random perturbation \( \delta \rho \) on the cloud density to make it moderately inhomogeneous, and \( \delta \rho/\rho \) is randomly distributed between \(-0.25 \) and \( 0.25 \). We set the characteristic
Table 1
Summary of Simulation Parameter

| Name     | \(\tau_{\text{IR}}\) | \(\tau_{\text{UV}}\) | \(F_{\text{UV}}\) \(|10^3|\) | \(F_{\text{IR}}\) \(|10^4|\) | \(T_{\text{eq}}\) \(|10^2|\) | \(D_c\) (pc) | Resolution (pc) |
|----------|------------------|-----------------|-----------------|-----------------|-----------------|----------------|-----------------|
| TLUV     | 0.1              | 12.658          | 2.0              | 3.075           | 0.411           | [4.110 $$\times 10^{-3}]$$ |                   |
| TSUV, L  | 0.01             | 1.266           | 2.0 $$\times 10^3$$ | 3.075           | 0.041           | [4.110 $$\times 10^{-3}]$$ |                   |
| TSUV, D  | 0.01             | 1.266           | 2.0 $$\times 10^3$$ | 3.075           | 0.411           | [4.110 $$\times 10^{-3}]$$ |                   |
| TSUV, DL | 0.001            | 0.127           | 2.0 $$\times 10^3$$ | 3.075           | 0.041           | [4.110 $$\times 10^{-3}]$$ |                   |
| TLIR, E  | 0.1              | 12.658          | 0.0              | 1.297           | 0.411           | [4.110 $$\times 10^{-3}]$$ |                   |
| TLIR, H  | 0.1              | 12.658          | 0.0              | 2.306           | 0.411           | [4.110 $$\times 10^{-3}]$$ |                   |
| TLMF, 0 | 0.1              | 12.658          | 1.0 $$\times 10^3$$ | 2.922           | 0.411           | [4.110 $$\times 10^{-3}]$$ |                   |
| TLMF, 1  | 0.1              | 12.658          | 2.0 $$\times 10^2$$ | 2.469           | 0.411           | [4.110 $$\times 10^{-3}]$$ |                   |
| TLUV, 3D | 0.1              | 12.658          | 2.0 $$\times 10^3$$ | 0.0             | 3.075           | [8.220 $$\times 10^{-3}]$$ |                   |
| TLUV, HR | 0.1              | 12.658          | 2.0 $$\times 10^3$$ | 0.0             | 3.075           | [2.055 $$\times 10^{-3}]$$ |                   |
| TLUV, LR | 0.1              | 12.658          | 2.0 $$\times 10^3$$ | 0.0             | 3.075           | [8.220 $$\times 10^{-3}]$$ |                   |
| TLUV, R  | 0.1              | 12.658          | 2.0 $$\times 10^3$$ | 0.0             | 3.075           | [4.110 $$\times 10^{-3}]$$ |                   |

Notes.
\(^a\) Flux in units erg s\(^{-1}\) cm\(^{-2}\).
\(^b\) \(T_{\text{eq}}\) in units of \(T_0 = 50\) K.
\(^c\) The number in the bracket gives the resolution at the finest level for each run in parsecs, the superscript outside the bracket gives the dimension of the corresponding simulation.

---

Table 2
Flying Distance and Final Velocity

| Name     | \(z\) (pc) | \(v_{\text{final}}\) (km s\(^{-1}\)) | \(t_{\text{final}}\) (yr) |
|----------|------------|-------------------------------------|---------------------------|
| TLUV     | 3.691      | 31.462                             | 2.502 $$\times 10^7$$     |
| TSUV, D  | 61.269     | 367.992                            | 3.314 $$\times 10^4$$     |
| TSUV, L  | 7.348      | 131.369                            | 1.099 $$\times 10^4$$     |
| TSUV, DL | 10.839     | 196.716                            | 1.071 $$\times 10^4$$     |
| TLIR, E  | 0.748      | 6.190                              | 2.502 $$\times 10^6$$     |
| TLIR, H  | 16.647     | 138.029                            | 2.502 $$\times 10^6$$     |
| TLMF, 10 | 20.147     | 203.361                            | 2.502 $$\times 10^6$$     |
| TLMF, 5  | 17.542     | 156.901                            | 2.502 $$\times 10^6$$     |
| TLMF, 1  | 16.560     | 139.304                            | 2.502 $$\times 10^6$$     |

Note.
\(^a\) For TSUV, D, TSUV, L, and TSUV, DL, which evolve over a sufficiently long period of time so that the cloud mass drops below 10% of its initial mass, \(t_{\text{final}}\) is defined as the time when the cloud mass drops to \(e^{-1.0}\) of the initial mass. For the rest of the simulations, \(t_{\text{final}}\) is the time that we stop the simulation. \(z\) and \(v_{\text{final}}\) are the flying distance and mean velocity corresponding to \(t_{\text{final}}\) respectively.

---

The length scale \(l_0\) to the initial cloud diameter \(D_c = 0.411\) pc, which corresponds to a column density of \(N_H = 7.566 \times 10^{22}\) cm\(^{-2}\). The corresponding initial optical depths to the IR and UV radiation are \(\tau_{\text{IR}} = 0.100\) and \(\tau_{\text{UV}} = 12.658\), respectively, so the cloud is optically thick to the UV radiation and optically thin to the IR emission. The cloud is embedded in the ISM with temperature \(T_{\text{iglu}} = 10^5\) K, and the background density is lower by \(10^{-5}\) to maintain pressure equilibrium.

With \(T_0\) and \(\rho_0\) defined, we introduce two dimensionless parameters:

\[
\Pi = \frac{a_T T_0^4}{\rho_0 v_0^2}, \quad \mathcal{C} = \frac{c}{v_0}.
\]

\(\Pi\) represents the ratio of the radiation pressure to the gas pressure, and \(\mathcal{C}\) represents the ratio of the light speed to the sound speed. With \(\Pi\) and \(\mathcal{C}\), the hydrodynamic equations and RT equation can be written in a dimensionless form (Jiang et al. 2014). The code solves the dimensionless versions of Equations (1), (3), and (10).

In the simulations, all of the hydrodynamic boundaries are set to outflow boundary conditions. Except for the lower \(x\) boundary, All radiation boundary conditions are set to outflow, which copies the radiation field in the last active zone, so that the radiation field in the calculation domain is isotropic for rays aligned with the positive \(x\)-direction even close to the vertical boundaries. At the lower \(x\) boundary, we impose a uniform and time-steady incoming radiation flux.

In the 2D simulations, the vertical (z) and horizontal (y or z) sizes of the simulation domain are \(40D_c\) and \(16D_c\), respectively, where \(D_c\) is the radius of the cloud. However, the 3D simulation has a smaller domain to reduce the computational cost. The size of the simulation domain is \(32D_c \times 6.4D_c \times 6.4D_c\), which spans \((-12D_c \times 20D_c) \times (-3.2D_c \times 3.2D_c) \times (-3.2D_c \times 3.2D_c)\), respectively. For all of the simulations, in order to better resolve the cloud, we apply a three-level static mesh refinement to the central region \((-5D_c \times 5D_c) \times (-2D_c \times 2D_c)\) (in the case of the 3D simulation, the refined central region in \(x, y, z\) is \((-5D_c \times 5D_c) \times (-2D_c \times 2D_c) \times (-2D_c \times 2D_c)\)). The resolution (cell size) at the most refined level is listed in Table 1. We also enforce a pressure floor \(P_{\text{floor}} = 10^{-5}\rho_0 v_0^2\) and a density floor \(\rho_{\text{floor}} = 10^{-5}\rho_0\). We use \(R = 10^{-2}\) and \(C_{\text{CFL}} = 0.4\) for all 2D runs except TLUV, R, in which we use \(R = 10^{-3}\). In TLUV, 3D, we adopted \(R = 10^{-2}\) and \(C_{\text{CFL}} = 0.3\).

We find it useful to define the cloud mass as the total mass of the cold gas, which is tracked by the dust and labeled by the passive scalar \(s\). Then, cloud mass is

\[
M_c \equiv \sum_i s_i \rho_i V_i
\]

where \(i\) runs over every grid cell in the simulation domain, and \(V_i\) is the volume of cell \(i\).

Initially (at \(t = 0\)), the passive scalar is set to be \(s = 1\) within the cloud, and \(s = 0\) in the background material. Hence, \(M_c\) is a representation of the mass of the dusty gas within the calculation domain. At later times, the cloud mass is the initial
mass $M_{c,0}$ minus the accumulative overheated gas mass:

$$ M_c(t) = M_{c,0} - \int M_{\text{loss}}(t) \, dt, $$

where $M_{\text{loss}}(t)$ represents the sum of all gas that has exited the domain or been lost to mixing with the hotter background gas. Note that the latter mechanism (mixing with hot gas) is the dominant loss channel in all simulations.

In order to focus on cloud evolution, we adopted a cloud-following frame approach, so the center of mass of the cloud remains fixed in the calculation domain.

In the cloud-following scheme, the x component of the mean velocity of dusty gas is computed at the end of every time step as

$$ \Delta v_{\text{mean}} = \frac{\int v_x \rho s \, dV}{\int \rho s \, dV}. $$

Then, $\Delta v_{\text{mean}}$ is subtracted from $v_x$ for every cell in the simulation domain. These boosts are then summed to keep track of the velocity $v_{\text{mean}}$ of the total cloud velocity after each time step. The hydrodynamics of the cloud is unaffected by these boosts due to the Galilean invariance of the underlying hydrodynamic equations. In contrast, the radiation equations are not Galilean invariant but Lorentz invariant. Hence, the radiation intensities differ at the second order in $v/c$ from the true Eulerian frame. For the calculations presented here, these discrepancies remain quite small and have an almost negligible impact on our results.

2.3. A Simple Model

A characteristic hydrodynamical timescale is set by the sound crossing time:

$$ t_0 = \frac{D_v}{v_0} \approx 6.255 \times 10^9 \left( \frac{50 \, \text{K}}{T_0} \right) \left( \frac{D_v}{0.411 \, \text{pc}} \right) \text{yr}. $$

There are several radiation timescales of interest. The first is the bulk acceleration time.

$$ \frac{\int \rho \, dv}{dt} = \int \frac{\rho_0}{c} F_{uv} e^{-\kappa_{uv} \rho_0 x^2} \, dx \, dA. $$

Integrating Equations (16) and assuming that the cloud is rigidly accelerated gives the average acceleration

$$ a = \frac{d\langle v \rangle}{dt} = \frac{\kappa_{uv} F_{uv} \rho_0}{c} \frac{1 - e^{-\tau_{uv}}}{\tau_{uv}}, $$

where $\tau_{uv} = \kappa_{uv} \rho_0 l_0$. When $\tau_{uv} > 1$, Equation (17) yields a characteristic acceleration rate $a \approx \kappa_{uv} F_{uv} / (\tau_{uv})$.

The UV radiation field heats the cold cloud while pushing on it. This heating is predominantly balanced by cooling via the IR radiation, so we estimate the approximate cloud characteristic temperature $T_{eq}$ by setting the UV absorption rate to balance the IR emission rate in the optically thin cloud:

$$ \kappa_{ir}(T_{eq}) T_{eq}^4 = \kappa_{uv} E_{uv} + \kappa_{ir}(T_{eq}) E_{ir}, $$

$$ = 2 \left( \frac{\kappa_{uv} F_{uv}}{c} + \frac{\kappa_{ir}(T_{eq}) F_{ir}}{c} \right). $$

The factor of 2 on the right-hand side comes from the moment integration of a half isotropic (isotropic for rays with $\mathbf{n} \cdot \hat{x} > 0$ and zero for rays with $\mathbf{n} \cdot \hat{x} < 0$) radiation field to get the energy density and flux in the $x$-direction. The resulting value is $T_{eq} = 153.8 \, \text{K}$ for the fiducial run.

The radiation field does not just accelerate and heat the cloud, it also acts to compress it. This is particularly true when the dusty cloud is opaque to the UV radiation and there is a significant radiation pressure gradient across the cloud. If we neglect the internal pressure support of the cloud, the time to crush it is simply determined by the relative acceleration of the cloud surface relative to the cloud center. With this assumption, we define a radiation crushing timescale as

$$ t_{\text{rad}} = \sqrt{\frac{D_v}{\Delta a_{uv}}} $$

where $\Delta a_{uv} = \kappa_{uv} F_{uv} (1 - e^{-\kappa_{uv} \rho_0 / 2}) / c$ represents the radiation acceleration difference between the outer radius of the cloud and its center due to the self-shielding of the UV flux.

3. Results

We preformed a series of 2D and 3D simulations to study various factors that impact on the cloud dynamics. We list the relevant parameters used in all simulations in Table 1. First, we report the fiducial run TLUV in Section 3.1, which is an opaque cloud accelerated by pure UV radiation flux. Parameters in this run were chosen with the expectation that they would provide favorable conditions for cloud survival. Next, we describe the impact of varying the optical depth in Section 3.2. Since the cloud responds differently to IR and UV radiation fluxes, we also report on the impact of varying the ratio of IR to UV flux in Section 3.3. We discuss the dependence of our results on dimensionality, resolution, and our choice for the reduced speed of light in Section 3.4.

3.1. UV Optically Thick Cloud

The TLUV run examines an optically thick cloud irradiated by a large UV radiation flux. This setup assumes that a large UV flux that would only be possible relatively close to a very compact, high-star-formation-rate region within a luminous galaxy. The column and gas densities are also quite large, so much so that the cloud would be unstable and gravitationally collapse if self-gravity were modeled here. A large column is chosen because we would like to explore the optical depth effects on the cloud. This combination might be expected to be favorable to cloud acceleration (large UV flux) and survival (significant self-shielding).

In this run, $\tau_{ir} = 0.1$, $\tau_{uv} = 12.7$. The incoming radiation flux from the lower $x$ boundary is pure UV flux. Figure 1 shows density and temperature snapshots from this simulation. The first row shows the density of both the hot and cold gas, and the second row shows only the temperature of the dusty gas ($s = 0$). Before the radiation flux comes into contact with the cloud, the cloud cools rapidly to the IR radiative equilibrium,
leading to an abrupt initial pressure drop. However, well before the cloud responds to the pressure mismatch with background medium, the incoming radiation field sweeps through the cloud. Immediately after the radiation reaches the cloud, the cloud temperature rises to around the estimated equilibrium temperature $T_{\text{eq}}$.

At early times ($t < 0.1 t_0$), the opaque cloud is compressed by the radiation pressure gradient within the cloud, which causes the side of the cloud facing the radiation field to be accelerated more strongly than the side opposite this face. Since the cloud is optically thick to the UV radiation, radiation only directly acts on the gas near the cloud surface. A dense distorted front is formed, and Rayleigh–Taylor-like instabilities grow at the interface between the hot and cold gas. The interior of the cloud is shielded from radiation and stays cold. As the radiation continues compressing the cloud, the gas pressure increases, with the cloud reaching its volume minimum near $t \sim 0.1 t_0$.

After this point, the gas pressure gradient counterbalances the radiation pressure, and the cloud partially re-expands and loses its initial spherical symmetry. In the re-expansion phase, the dense core of the cloud remains cold ($T \sim T_{\text{eq}}$) and is stretched slightly, primarily along the direction of motion. The lower-density envelope of gas becomes turbulent and filamentary. Although the center of mass of the dusty gas remains fixed on the grid, the boosts associated with the cloud-following scheme result in a significant velocity in the background medium toward the bottom $x$ boundary, roughly at the value of $v_{\text{mean}}$ in Figure 2. The combination of the large relative velocity of the background flow and the radiation pressure from the UV drive Kelvin–Helmholz-like instabilities that facilitate mixing with the hotter background gas. On any single time step, a small fraction of the gas is heated above the assumed destruction temperature. The detailed morphology of low-density gas at late times is sensitive to the assumed initial condition, which we tested by considering different random

**Figure 1.** Simulation snapshots from TLUV. Top panels: density snapshots of both cold and hot gas. Lower panels: temperature of the cold dusty gas; the background medium is masked by black. The maximum temperature in the color bar corresponds to the temperature at which we set the passive scalar to zero, representing the overheating of cold gas. $t_0 = 6.255 \times 10^5$ yr, $D_c = 0.411$ pc, $\rho_0 = 10^{-19}$ g cm$^{-3}$, and $T_0 = 50$ K.

The Astrophysical Journal, 893:50 (15pp), 2020 April 10 Huang, Davis, & Zhang
perturbations, but the qualitative picture of compression, re-expansion, and mixing outlined here was qualitatively similar in all runs.

In Figure 2, the black solid line shows the evolution of the cloud mean velocity $v_{\text{mean}}$ (top panel), velocity dispersion $\sigma_v$ (middle panel), and cloud mass $M_c$ (bottom panel) evolution of TLUV (black), TSUV_D (green), TSUV_L (blue), and TSUV_DL (orange). In the top panel, the dashed line with the same color is the time integration of $a$ (Equation (17)) of each run, except for TLUV and TLUV_L, which have the same $a$. In the bottom panel, the solid lines are $M_c$, and the dashed line with the same color is $M_c(t)$ of each run. We ended the simulations when $M_c$ and $M_c(t)$ started to diverge, meaning that the cold gas exiting the simulation box starts to effect total mass loss. $v_0 = 0.642 \text{ km s}^{-1}$.

![Figure 2](image)

**Figure 2.** Mean velocity $\Delta v_{\text{mean}}$ (top panel), velocity dispersion $\sigma_v$ (middle panel), and cloud mass $M_c$ (bottom panel) evolution of TLUV (black), TSUV_D (green), TSUV_L (blue), and TSUV_DL (orange). In the top panel, the dashed line with the same color is the time integration of $a$ (Equation (17)) of each run, except for TLUV and TLUV_L, which have the same $a$. In the bottom panel, the solid lines are $M_c$, and the dashed line with the same color is $M_c(t)$ of each run. We ended the simulations when $M_c$ and $M_c(t)$ started to diverge, meaning that the cold gas exiting the simulation box starts to effect total mass loss. $v_0 = 0.642 \text{ km s}^{-1}$.

The average temperature is always well below the dust destruction temperature $T_{\text{eq}}$ in the compression phase, and gradually rises toward $T \approx T_{\text{eq}}$ in the re-expansion phase, but never quite reaches it. This is primarily a problem with our $T_{\text{eq}}$ estimate, which assumes the energy density and flux related by $E_{\text{uv}} = 2F_{\text{uv}}/c$. The energy density is actually lower than this due primarily to the attenuation of the incoming UV and (minor) differences in the assumed angular distribution, leading to a slightly smaller $T_{\text{eq}}$ than our estimate predicts. The rise in temperature can be attributed primarily to more gas being exposed to the incident radiation field as the cloud re-expands but mixing with the hotter background plays a role as well.

The average temperature is always well below the dust destruction temperature, indicating that only a small fraction of the dusty gas in the envelope is mixing with the background on any time step. This suggests there is a continuous flux of cooler, higher density gas leaving the compressed cloud core for the envelope, mixing with the background, and being destroyed. As discussed in Section 4.1 this process happens via mixing on the radiation hydrodynamic timescales. Efficient IR cooling prevents the radiation from simply overheating the cloud on the radiation crossing time.

3.2. UV Optically Thin Clouds

As noted above, the TLUV run corresponds to a rather large cloud mass and column. Here, we consider UV optically thin runs, with lower cloud columns that might be more typical of outflowing gas in star-forming environments. We studied two runs with optical depth of about unity, both with $\tau_{\text{uv}} = 1.266$, a factor of 10 lower than TLUV. The UV optical depth $\tau_{\text{uv}} = \kappa_{\text{uv}} n_{\text{H}_2} D_c$, can be made smaller by either reducing the radius or lowering the density. The cloud with lower density is in TSUV_D, the cloud with smaller radius is in TSUV_L. For a third run, TSUV_DL, we reduce both the density and length, giving $\tau_{\text{uv}} = 0.127$. In all cases, the optical depth to IR radiation $\tau_{\text{IR}}$ is reduced by the same factors.
Figure 2 compares the cloud bulk motion of optically thin clouds to the optically thick cloud in TLUV. In both TSUV_L (blue) and TSUV_D (green), estimated acceleration in Equation (17) (shown as a dashed curve) is about 10 times larger than TLUV. Both optically thin runs roughly follow this curve, falling slightly below during the re-expansion stages. The compression and re-expansion is qualitatively similar to what is seen in the TLUV run, leading to similar evolution of velocity dispersion and cloud mass, albeit on somewhat different timescales. For example, cloud mass is already substantially depleted by $t = 10^5$ yr in the TSUV_L and TSUV_DL runs. The TSUV_DL run is accelerated by a greater amount consistent with the expectations from Equation (17), and the velocity dispersion tends to be lower but the mass loss initially follows the TSUV_L run and is destroyed on a comparable timescale. The mass loss in the TSUV_D run follows a similar evolution to that in the TLUV run, but we run it for longer because it takes longer for the outflow through the domain boundaries to dominate dust mass loss. We see that the runs with shorter length scale are disrupted much faster, consistent with the shorter cloud crushing timescale implied by Equation (19). Note that the drops in velocity dispersion near the end of the runs follow in part from no longer being able to track mass leaving the domain.

Figure 4 compares the density snapshots of the optically thin clouds. The morphology and evolution of TSUV_L and TSUV_D are qualitatively similar. Due to the smaller optical depth, the radiation is more uniformly distributed than in the TLUV run. The clouds do not show a dense front at the cloud-radiation interface like in TLUV, where the radiation is absorbed. Hence, the clouds are more uniformly compressed by radiation pressure at early times. As the clouds re-expand, the core of the dense gas is somewhat stretched along the direction of motion. As in TLUV, lower-density gas pushed by both radiation pressure and the interaction with the background gas leads to a turbulent, filamentary structure. The low-density regions ejected by the cloud mix with a hot background, gradually heat the gas to the destruction temperature. The TSUV_DL run follows a similar evolution at early times, but there is much less overall compression because the radiation field is much more uniform. In fact, the average density in the cloud drops as the outer layers expand. Nevertheless, the subsequent evolution is qualitatively similar to the re-expansion phases of the other runs. Shear at the interface with the background ISM disrupts the cloud and drives mixing with the hotter background and heats the cold gas to $T > 1500\,K$.

Figure 3 also shows the difference in their average density $\bar{\rho}$ and temperature $T$ for the optically thin runs. As already noted, the TSUV_L run is compressed faster than are the runs with larger length scales, but it has a lower maximum density. The TSUV_D run is also compressed slightly faster than is the TLUV run, but it reaches a much higher maximum average density. Note that TSUV_D has a lower cloud and background density, so the density and temperature contrast between the cloud and background is the same as those for TLUV. As noted above, the TSUV_DL run does not show the compression and re-expansion behavior that is seen in the more optically thick runs. The average density drops continuously. During the initial compression phase, the TSUV_L and TSUV_D runs remain below the $T_{eq}$ estimate due to a lower $E_{uv}$ in the cloud core, but they rise to near this value at later times as the optically thin assumption underlying this estimate holds better in these runs than in the TLUV runs. The evolution is roughly similar to the optically thin simulations in Proga et al. (2014) but with larger relative velocity between the cloud and background.

The differences in evolution between the optically thick and optical depth unity runs can be attributed primarily to the gas pressure distribution within the clouds. Initially, both clouds are in pressure equilibrium with the background medium. After radiation has swept through the cloud, different parts of the cloud experience different radiation forces. The radiation pressure gradients are modest in the optically thin runs, but self-shielding leads to strong gradients in the UV optically thick run. These radiation pressure gradients compress the cloud until a comparable gas pressure gradient develops to oppose it. Figure 5 shows the pressure snapshots of TSUV_L, TSUV_D, TSUV_DL, and TLUV at the same compression stage. TSUV_L and TSUV_D are at optical depths of unity, so the gas pressure gradient is relatively small. In contrast, the gas pressure is strongly enhanced near the surface in the optically thick TLUV runs (right panels). The outward pressure gradient supports the cloud and slows down the compression. TSUV_DL is even more optically thin and shows almost no pressure gradient at early times.
Initially, these gas pressure effects are modest, and our estimate of the cloud crushing time $t_{\text{rad}}$ in Equation (19) yields a good order-of-magnitude estimate of the time for both the optically thick and optical depth unity clouds to reach their maximum average density. However, if we look more quantitatively, we can see the impact of the gas pressure gradient force, which is not accounted for Equation (19). Our estimates of $t_{\text{rad}}$ are shown as the vertical dashed lines in Figure 3. They provide better estimates of the time when the cloud reaches peak density in the two optical depth unity runs, but the maximum compression of the optically thick cloud is slightly delayed due to the resistance from the gas pressure gradient. Note that the compression is somewhat faster in the TSUV_L run because $\tau_D = 0.1$, $D_c = 0.1D$, Characteristic pressure $P_0 = \rho_0 v_0^2$, $R_0 = 4.127 \times 10^{-10}$ dyne cm$^{-2}$ for TSUV_L and TLUV, and $P_0 = 4.127 \times 10^{-11}$ dyne cm$^{-2}$ for TSUV_D and TSUV_DL.

Since the temperature is similar in both optical depth unity runs, the lower initial density in the TSUV_D run means that it has a lower initial gas pressure than the TSUV_L run. Since the cloud reaches its peak density when the gas pressure gradient becomes large enough to support the cloud against radiative compression, its largest peak average density is $\bar{\rho}_{\text{peak}}/\rho_{\text{init}}$ (green solid line in Figure 3).

3.3. Acceleration with Both IR and UV Irradiation

In contrast to UV radiation flux, cloud acceleration with IR radiation is generally gentler due to the smaller optical depth and the fact that IR radiation acts both to compress the cloud (incident radiation) but also provides support against compression (re-radiated IR). Previous work has suggested that clouds accelerated solely by IR radiation might survive longer than clouds entrained in a hot outflow (Zhang et al. 2018). Hence, we have performed a number of runs with an incident IR flux to compare with the pure UV results discussed above. We first consider two purely IR runs: (1) TLIR_E with an incident IR flux equal to the UV runs above, and (2) TLIR_H, which has a flux a factor of 10 larger. For both cases, we set the initial IR optical depth $\tau_{\text{IR}} = 0.1$.

Figure 6 shows the cloud mean velocity $v_{\text{mean}}$ (top row), velocity dispersion $\sigma_v$ (middle row), and cold gas mass $M_c$ (bottom row) for TLIR_E (green) and TLIR_H (red), with TLUV (black) included for comparison. TLIR_E is accelerated much more slowly than TLUV due to its smaller opacity, which gives rise to an acceleration $a_{\text{IR}} \approx \kappa_{\text{IR}} F_{\text{IR}}/c < a_{\text{UV}}$ when $F_{\text{IR}} = F_{\text{UV}}$. Since the IR-driven cloud is not compressed significantly, the cloud is more weakly disrupted and retains its initial structure longer, leading to a lower velocity dispersion. The mean acceleration is nearly constant, giving rise to a nearly linear velocity profile. A similar evolution is seen for TLIR_H, but the factor of 10 increase in $F_{\text{IR}}$ compensates for the lower opacity and ultimately leads to a more rapid acceleration than is seen for TLUV.

Figure 7 shows the cold gas average density (upper panel) and temperature (lower panel) for TLIR_H (black) and TLIR_E (orange). In contrast to the UV runs, both the mean density and temperature remain relatively constant for these runs. In fact, the average density shows a slight drop as the radiation pressure associated with the re-emitted IR leads the cloud becoming weakly over-pressured relative to the background and expanding slightly. There is a brief initial transient when the radiation sweeps across the cloud and heats it to the equilibrium temperature, where it remains for the rest of the evolution. Overall, our results are in good agreement with the large-scale optically thin cloud simulations performed by Zhang et al. (2018).

Given the disparate evolution histories and survival times in the IR-only and UV-only runs presented thus far, it is natural to ask how a combination of UV and IR driving affects cloud evolution. For highly star-forming galaxies, such as ULIRGs, the UV usually represents a small fraction $\lesssim$1% of the total observed emission. Most of this radiation is thought to be originally emitted in the optical and UV by stars and then reprocessed in the IR due to the large dust optical depths along...
most lines of sight. However, it is plausible that the UV will have been less attenuated in the location where the outflows are launched, motivating an exploration of different ratios of UV-to-IR flux.

We consider three simulations all with the same incident IR flux, which is equivalent to TLIR_H. These simulations also have an incident UV flux corresponding to 1% (TLMF_1), 5% (TLMF_5), and 10% (TLMF_10) of the IR flux. The parameters are listed in Table 1. The cloud is optically thick to UV radiation ($\tau = 12.7$) and optically thin to IR radiation ($\tau_{IR} = 0.1$).

As with UV-only runs, the UV provides a differential radiation force that acts to first compress the cloud until the gas pressure rises and drives re-expansion. This density evolution is seen most clearly for TLMF_5 and TLMF_10 in Figure 7, but is only modest for TLMF_1, which is similar to the IR-only runs. Density snapshots shortly after maximum compression (at $0.15t_0$) are shown in Figure 8, where one can see the increasing degree of disruption as the UV fraction increases. It’s also interesting to compare the average cold gas density of TLMF_10 and TLUV in Figure 7. Despite the same incident UV flux, TLMF_10 has a lower peak density than TLUV because of the support from re-emitted IR radiation against UV radiation. As the cloud is heated, its average temperature increases, TLMF_1 approaches the optically thin equilibrium radiation. As the cloud is heated, its average temperature increases, TLMF_5 approaches the optically thin equilibrium radiation. As the cloud is heated, its average temperature increases, TLMF_10 approaches the optically thin equilibrium radiation.

Figure 9 shows the mean velocity $\Delta v_{mean}$ (top panel), velocity dispersion $\sigma_v$ (middle panel), and cloud mass $M_c$ (bottom panel) evolution of TLMF_10 (red), TLMF_5 (blue), TLMF_1 (green), and TLIR_H (black). $v_0 = 0.642 \text{ km s}^{-1}$, and $t_0 = 6.255 \times 10^5 \text{ yr}$. Here, $\rho_0 = 10^{-19} \text{ g cm}^{-3}$.

The overall impression is that the acceleration is enhanced if there is enough UV flux so that the radiation force from the UV matches or exceeds the IR radiation force. However, the survival of dusty gas is sensitive to the relative contributions of UV and IR radiation that drives it. Since the temperature of the majority of the gas remains close to the equilibrium temperature (see Table 1 and Equation (18)), it is not a matter of the

---

**Figure 7.** Average density (top panel) and temperature (bottom panel) of cold gas for multifrequency runs TLMF_10 (red), TLMF_5 (blue), and TLMF_1 (green) and IR radiation runs TLIR_H (black) and TLIR_E (orange). $\rho_{\text{init}} = \rho_0 = 10^{-19} \text{ g cm}^{-3}$, $t_0 = 50 \text{ K}$.

**Figure 8.** Density snapshots of TLMF_10 (left panel), TLMF_5 (middle panel), and TLMF_1 (right panel) at $t = 0.15t_0$. Here, $t_0 = 6.255 \times 10^5 \text{ yr}$.

**Figure 9.** Mean velocity $\Delta v_{mean}$ (top panel), velocity dispersion $\sigma_v$ (middle panel), and cloud mass $M_c$ (bottom panel) evolution of TLMF_10 (red), TLMF_5 (blue), TLMF_1 (green), and TLIR_H (black). $v_0 = 0.642 \text{ km s}^{-1}$, and $t_0 = 6.255 \times 10^5 \text{ yr}$.
3.4. Dimensionality, Resolution, and Reduced Speed of Light

The hydrodynamic interactions that drive mixing and cloud destruction are potentially sensitive to resolution. We considered the effects of resolution by rerunning our fiducial run at two additional resolutions. The TLUV_HR and TLUV_LR runs are performed at resolutions that are a factor of two higher and lower, respectively, in both dimensions relative to the TLUV_run (see Table 1). Figure 10 shows the bulk motion and mass evolution of the cloud for different resolutions. The evolutions of the mean density and temperature are shown in Figure 12.

The motion and mass evolution of the runs are all qualitatively similar to each other, although there are modest deviations later in the evolution, when the nonlinear effects start to dominate. These later time deviations are at about the same level that we see when changing the random initial perturbations on the cloud density. Slightly more sensitivity is seen in the evolution of the average density, where the maximum average density reached scales with resolution, suggesting that the core cloud structure is not yet resolved at the peak. However, almost all runs asymptote to similar values of density at late times and show little variation in temperature evolution. Hence, our results do not seem to be substantially impacted by resolution for the conditions considered here.

All of the simulations presented above were 2D, but the hydrodynamic effects that lead to mixing with the background might depend on dimensionality, so we also performed a 3D simulation. Since 3D runs are considerably more expensive, we only carry out one run (TLUV_3D) to see how well our 2D results generalize to 3D. Due to the increased computational cost, the resolution of TLUV_3D is chosen to be equivalent to that of the TLUV_LR run. Comparing these runs in Figure 10, we find that the acceleration and survival times are rather similar for both runs.

Figure 11 compares a 2D slice from TLUV_3D with a density snapshot at the same time in the TLUV_LR run. The hot background medium is masked by black. The hot background medium is masked by black. Figure 12 shows the bulk motion and mass evolution of the cloud for different resolutions. The evolutions of the mean density and temperature are shown in Figure 12.

The motion and mass evolution of the runs are all qualitatively similar to each other, although there are modest deviations later in the evolution, when the nonlinear effects start to dominate. These later time deviations are at about the same level that we see when changing the random initial perturbations on the cloud density. Slightly more sensitivity is seen in the evolution of the average density, where the maximum average density reached scales with resolution, suggesting that the core cloud structure is not yet resolved at the peak. However, almost all runs asymptote to similar values of density at late times and show little variation in temperature evolution. Hence, our results do not seem to be substantially impacted by resolution for the conditions considered here.

Figure 10. Mean velocity $\Delta v_{mean}$ (top panel), velocity dispersion $\sigma_v$ (middle panel), and cloud mass $M_c$ (bottom panel) evolution of TLUV_3D (green), TLUV (black), TLUV_LR (blue), TLUV_HR (red), and TLUV_R (orange). In TLUV_R, the radiation flux travels 10 times slower than in other runs because of the lower reduction factor. $v_0 \approx 0.64 \text{ km s}^{-1}$, $t_0 \approx 6.255 \times 10^5 \text{ yr}$.

Figure 11. Dust density snapshots of TLUV_LR (left panel) and TLUV_3D (right panel) at $t = 0.06t_0$. The hot background medium is masked by black. $t_0 = 6.255 \times 10^5 \text{ yr}$, $\rho_0 = 10^{-19} \text{ g cm}^{-3}$.

Figure 12. Average density (upper panel) and temperature (lower panel) of cold gas in TLUV_3D (green), TLUV (black), TLUV_LR (blue), TLUV_HR (red), and TLUV_R (orange). $\rho_{0,\text{hot}} = \rho_0 = 10^{-19} \text{ g cm}^{-3}$, $T_0 = 50 \text{ K}$. In TLUV_R, the radiation flux travels 10 times slower than in other runs because of the smaller reduction factor. We moved the curves of TLUV_R 0.01 earlier in order to compare the cloud dynamics with other runs.
homologously in three rather than two dimensions. However, as with resolution, the simulations asymptote to similar densities at later times, and we conclude that dimensionality has relatively little effect on the cloud acceleration or survival time for this setup.

Finally, we also test the effect of the speed-of-light reduction factor $R$ on cloud dynamics. In TLUV$_R$, we chose a reduction factor of $R = 10^{-3}$. In other words, the speed of light is 10 times smaller than that in the TLUV run, allowing for time steps that are 10 times larger. As a result, it takes 10 times longer for the radiation from the lower $x$ boundary to reach the cloud. Once this offset is accounted for, the cloud bulk motion, velocity dispersion, mass (Figure 10), and mean density and temperature evolution (Figure 12) are all qualitatively similar to TLUV. The primary difference is that the peak average density is lower in this run, which is to be expected since the sound speed is closer to the reduced speed of light. This allows for more gradual adjustments in the gas pressure as the radiation sweeps across the cloud. Nevertheless, the density asymptotes to a similar mean cloud density at later times. Hence, we do not believe our simulations are sensitive to our choice for $R$.

4. Discussion

4.1. Destruction Mechanism for Cold Gas

Our primary interest in this problem is assessing whether radiative acceleration could play an important role in accelerating the outflows observed in molecular and atomic transitions. Therefore, an important constraint is that the gas cannot be too hot for the observed transitions to be present. Since the optimal temperature ranges for different species can vary significantly, there is no single temperature cutoff that describes all transitions. We have utilized the presence of dusty gas $s > 0$ as the criterion for survival. Since we have a adopted 1500 K as our dust destruction/decoupling temperature, this can be thought of as a proxy for molecular gas. We do not believe our results are significantly sensitive to this choice of temperature because the bulk of the cloud stays close to the radiative equilibrium temperature $T \lesssim 100$ K, and only a modest fraction has temperatures that are significantly higher than this. The dusty gas (defined by $s > 0$) at a temperature significantly higher than the equilibrium temperature is being rapidly heated by mixing with the background gas to temperatures near the assumed background temperature ($T \gtrsim 10^5$ K). For this reason, test runs with a lower destruction/decoupling temperature (500 K or 1000 K instead of 1500 K) are not significantly different, because the gas reaching 500 K continues heating and quickly exceeds 1500 K shortly after. Hence, our results should not be sensitive to the assumed destruction temperature as long as this temperature is well below the background temperature and well above the equilibrium temperature.

As discussed in Section 3, the process of cloud disruption and heating is primarily a radiation hydrodynamical process rather than simply an RT process. In other words, the cloud is not simply heated to a high temperature by the UV flux. The efficient IR dust cooling allows the majority of gas to remain close to the radiative equilibrium temperature. Instead, the radiation pressure forces drive compression and re-expansion of the gas in optically thick clouds. In the optically thin limit, the compression and re-expansion phase is much more subdued, but the dynamics still leads to mixing of the outer layers. In all simulations, a significant fraction of lower-density dusty gas in the outer envelope of the cloud is heated by this mixing with the hotter background ISM resulting in the destruction of most of the cloud on relatively short timescales.

Figure 13 shows the density distribution of hot and cold gas in the cloud ($s = 0$) for the TLUV run at three different times. The blue solid lines are the density distribution of cold gas (left $y$-axis labels). The red solid lines are the density distribution of the hot gas (right $y$-axis labels), which we defined as gases with temperatures higher than 95% of the assumed dust destruction temperature. The top panel shows before compression, the absence of a red solid line means at $t = 0.04t_0$, there is no hot dusty gas. The middle panel shows before the volume minimum, where the mean cloud density is near its extremum. The hot dusty gas (red) has a density about two orders of magnitude smaller than the majority of the cold dusty gas (blue). The bottom panel shows the density distribution during the re-expansion, with hot gas still corresponding to lower densities than most of the cold gas. In effect, there is a continuous flux of cold dense gas toward lower densities due to the expansion and mixing that is heated and incorporated into the background ISM.

After the cloud begins to be disrupted, the implied rapid mixing with the background ISM is consistent with expectations from purely hydrodynamic models, driven by Kelvin–Helmholtz instabilities (Begelman & Fabian 1990). We assume a mixing timescale $t_{\text{mix}} \sim t_{\text{KH}} \sim (\rho_c/\rho_h)^{1/2} t_{\text{edd}}$, where $\rho_c$ is the density of the cold material, $\rho_h$ is the density of hot material, and $t_{\text{edd}}$ is the timescale for eddies to cross the mixing layer. We estimate $t_{\text{edd}} \sim 0.1t_0/\sigma_v$ with $\sigma_v \sim 10–50$ across different runs. For the relatively low-density envelope of dusty gas, we have ratios $\rho_c/\rho_h \sim 10^2–10^3$, yielding $t_{\text{mix}} \approx 0.02t_0–0.3t_0$. Therefore, it is not surprising that gas driven to lower densities by hydrodynamic processes in the cloud surface can rapidly mix with the background. However, if we apply the same argument to the initial cloud with a density ratio $10^5$ times larger than the background, the timescale is much longer,
indicating that radiative forces play an important role in the disruption.

4.2. Cloud Survival Time

A number of numerical studies of purely hydrodynamic entrainment (Scannapieco & Brüggen 2015; Brüggen & Scannapieco 2016; Zhang et al. 2018) concluded that entrainment in a hot wind is unlikely to accelerate the cloud to the observed speed before the cloud is shredded. However, Zhang et al. (2018) found that clouds accelerated by an IR radiation field can survive longer if a sufficiently larger IR flux is available to accelerate them. Our focus here was to consider the degree to which the addition or substitution of UV radiation impacts this conclusion.

Following Zhang et al. (2018), we define the cloud surviving time as the time when the cloud loses half of its initial mass. Our definition of cloud mass follows from Equation (12), which includes mass loss both from mixing and advection out of the domain. In Figure 14, the solid lines correspond to $M_{c,int}$, and the dashed line with same color excludes mass loss associated with advection through the domain boundary. In other words, it assumes (conservatively) that cold gas advected through the boundary remains cold, and only gas overheated within the domain is accounted for. Since we stop the simulations when the mass loss out of the domain begins to become significant, almost all of the mass loss shown in the figures is due to mixing and overheating within the simulation domain.

A potentially important characteristic timescale for estimating the cloud survival time is the radiation crushing time $t_{rad}$ from Equation (19). For runs with $\tau_{uv} \geq 1$, the mass loss occurs primarily after the maximum compression, when the cloud starts to re-expand and the timescale for re-expansion is comparable to or slightly longer than $t_{rad}$. Hence, for this regime, it provides an approximate estimate of the survival time but the dependence on $\tau_{uv}$ is not borne out. It implies that the most optically thick clouds will be crushed the fastest, but this is not what we found. The difference arises because $t_{rad}$ neglects the impact of the pressure support from gas and re-radiated IR radiation.

Since the UV radiation acts to compress the cloud without providing pressure support, one might expect the optically thick cloud accelerated by UV radiation alone to have a shorter surviving time than the one supported by IR alone, and this is consistent with our results. Figure 14 compares the cloud-mass evolution of the UV-driven runs that are optically thick (TLUV) to optically thin (TSUV_D, TSUV_L, and TSUV_DL) with the IR-driven run (TLIR_H). The cloud mass drops significantly faster when accelerated by UV radiation. For clouds with optical depths of the order of unity or smaller, one expects the higher opacity of UV radiation to result in a more rapid acceleration (see Equation (17)). However, the significantly shorter survival time still limits the velocity to less than is typically inferred from observations.

With the exception of the pure IR runs (which are run for the same time as the fiducial run), we end all of the simulations when the outflow out of the simulation domain starts to become the dominant mass-loss mechanism. At this point, TLUV only has a speed $v_{mean} \approx 31.5 \text{ km s}^{-1}$, while $M_{c}/M_{c,int} \approx 60\%$, which is much smaller than the observed velocities, which are 100 s of km s$^{-1}$ (Oosterloo et al. 2017; Krieger et al. 2019). Integrating the mean velocity over time, we can estimate the "flying" distance $z \sim \int v_{mean} dt \approx 3.7 \text{ pc}$, which is small compared to the typical size of the star-forming region. Therefore, we do not expect optically thick UV-driven clouds to survive long enough to explain observed outflows. TSUV_L has larger acceleration and, therefore, reaches larger velocities of $\sim 131.4 \text{ km s}^{-1}$ by $t_{final}$, but flying distance remains small due to the short survival time. The most promising UV run is TSUV_D, which survives longer than TSUV_L and reaches velocities of $\sim 368.0 \text{ km s}^{-1}$ when the mass drops to $\sim 36\%$ of its initial mass.

In contrast, IR radiation flux accelerates clouds that are optically thin to IR for several dynamical timescales with relatively little disruption (Zhang et al. 2018). Incident radiation momentum is converted efficiently into cloud momentum. Although there is still some turbulent motion in the cloud outer layers and associated mixing with the background, it is much less than in the UV, and the cloud survival time is much longer.

4.3. Effects of Multiband Irradiation

Since a cloud absorbs both UV and IR radiation but only re-emits IR radiation (Equation (7)), the UV and IR radiation interact very differently with the clouds. UV radiation tends to accelerate the opaque cloud faster, but also compresses the cloud and eventually drives greater mixing with the hot background ISM. In contrast, the IR radiation flux accelerates the cloud more uniformly without significant compression, leading mixing to occur on much longer timescales. A cloud’s interaction with a mixture of IR and UV flux is somewhat more complicated. Figure 15 shows the cloud-mass evolution from the multiband irradiating flux runs. TLMF_10 has the same UV flux as TLUV (black) and the same IR flux as TLIR_H (red). Despite the presence of extra IR radiation flux that is of an order of magnitude higher than the UV flux, the cold gas mass

![Figure 14](image-url) Figure 14. Cloud-mass evolution for different runs. The horizontal gray dashed line represents when the cloud mass is half of initial mass, corresponding to the cloud surviving time. The black lines show TLUV. The blue lines show TSUV_L, the green lines show TSUV_D, and the orange lines show TSUV_DL. The red lines show the mass evolution of the optically thick cloud in TLIR_H, which is accelerated by pure IR radiation. For each color, the solid line is $M_{c}(t)$ (Equation (13)).
The primary goal of this work is to examine the relative role of UV and IR radiation pressure in accelerating outflows of cold molecular gas that are observed in star-forming galaxies. Since we focus on this mechanism, we implicitly ignore other possibilities such as entrainment in hot outflows or acceleration due to cosmic-ray pressure (e.g., McCourt et al. 2015; Scannapieco & Brüggen 2015; Zhang et al. 2018; Wiener et al. 2019). In principle, these other acceleration mechanisms may all act in concert to drive outflows, or radiation may be an entirely subdominant process. Our primary motivation for neglecting other acceleration mechanisms is that it allows us to focus on and better understand the radiation hydrodynamics; but given the uncertainties, we view radiation-dominated acceleration as a physically plausible limit.

Due to the expense of solving the RT equations, even the mostly 2D simulations presented here are relatively computationally expensive. This requires us to make trade-offs in our modeling. Our simulations focus on the radiation hydrodynamics of the acceleration process but improve upon earlier work by studying the interaction of dusty gas with a multiband irradiating flux under differing assumptions about the optical depths and UV and IR fluxes. However, we simplify or neglect some of the more complex physics that may be relevant to realistic outflows. Future work could benefit from including the overall galaxy gravitational potential and any self-gravity of the gas, studying the effects of the magnetic fields (McCourt et al. 2015), modeling dilution of the radiation field far from the original source, including additional complexity in the background ISM, a detailed treatment of photoionization (e.g., Namekata et al. 2014), modeling of conduction (Brüggen & Scannapieco 2016), or the process of cloud formation (e.g., Proga & Waters 2015; Waters & Proga 2016). We also utilized a simplified prescription for the dust opacity (Semenov et al. 2003) and adopted a simple temperature criterion to determine when the dust is destroyed or decoupled (Krumholz & Thompson 2013) from the gas. Future work may benefit from more elaborate treatments of the dust, including its coupling to gas, destruction mechanisms, and opacity.

5. Conclusion

We consider the effect of UV radiation pressure on the acceleration of cold, dusty gas. In contrast to earlier work that focused on IR radiation alone, we find that replacing the IR with UV radiation or including a large fraction of UV radiation is generally detrimental to cloud survival. This is due to the UV radiation pressure distorting and compressing the cloud, driving mixing with the hotter background ISM, with mixing ultimately leading to overheating and dust destruction. In contrast, simulations dominated by IR radiation are more robust, with longer survival times, in agreement with earlier work (Zhang et al. 2018). We attribute this difference to IR radiation’s role in both accelerating the cloud but also in providing an internal radiation pressure due to dust reemission that maintains a more uniform cloud structure and limits mixing.

We also consider the impact of the optical depth on cloud dynamics. All simulations considered here are optically thin to the IR but range from optically thin ($\tau_{\text{uv}} = 0.13$) to optically thick ($\tau_{\text{uv}} = 13$) in the UV. Generally speaking, decreases in the UV optical depth of the cloud lead to faster disruption evolution is similar to TLUV. Such behavior also reflects the fact that the overheating of cold gas is dominated by mixing with the hot background medium, which is triggered by the cloud re-expansion after interacting with the UV radiation. Since the UV fluxes are the same, the surviving timescale is similar. However, the additional IR radiation provides a larger acceleration, and the cloud in TLMF_10 reaches a higher final velocity. By the time we stopped the simulation, TLMF_10 is accelerated to \( \sim 203.4 \text{ km s}^{-1} \) and has moved 20.1 pc from its start position, and about 50% of cold gas in the cloud has been heated over 1500 K.

Lowering the UV flux fraction in the TLMF_5 and TLMF_1 runs produces a mass evolution increasingly similar to the pure IR irradiation case. Note that the difference in mass and average density (Figure 7) evolution between 5% and 1% of UV radiation flux is somewhat sharp. Comparing the different UV fraction runs in Figure 9 shows that increasing the UV radiation from 1% to 5% does not significantly increase the bulk acceleration because of the small absolute value of the UV flux. But the effect on the compression of the cloud is much more significant. Hence, even a modest UV fraction can disrupt the cloud without significantly improving the acceleration. However, the similarity between the TLUV and TLMF_10 runs also suggests that it is not simply the UV fraction that matters since these runs have UV fractions of 100% and 9% of the total (IR and UV) flux, respectively. Our results suggest that if the UV flux and optical depth are large enough to produce significant differential acceleration, compression of the cloud will drive re-expansion and disrupt the cloud in a manner that will significantly enhance mixing with the hot background flow. This would suggest that the most optically thick star-forming environments, such as ULIRGs where the vast majority of stellar light is reprocessed into the IR, may be the most efficient locations for driving molecular outflows if radiation pressure dominates.

4.4. Model Uncertainties and Approximations

The Astrophysical Journal, 893:50 (15pp), 2020 April 10

Huang, Davis, & Zhang
radiation pressure is most likely to play a role in driving star-forming galaxies are the environments where considerations combined with the need for large radiation pressure acceleration will be most relevant in highly obscured star-forming galaxies where the UV fractions are low. In contrast, disruption and mixing will likely tend to destroy clouds in more UV-dominated starburst galaxies. These considerations combined with the need for large radiation fluxes (Zhang et al. 2018) suggest that ULIRGs and high-redshift star-forming galaxies are the environments where radiation pressure is most likely to play a role in driving outflows of cold molecular gas.

We thank the referee for a helpful and detailed report that significantly improved this work. We also thank Evan Scannapieco, Justin Spilker, and Eve Ostriker for helpful conversations as well as Yan-Fei Jiang and Jim Stone for their contributions to the Athena++ radiation module used in this work. This work used the computational resources provided by the Advanced Research Computing Services (ARCS) at the University of Virginia. We also used the Extreme Science and Engineering Discovery Environment (XSEDE), which is supported by National Science Foundation (NSF) grant No. ACI-1053575. This work was supported by the NSF under grant AST-1616171.

Facilities: Rivanna.
Software: Athena++, VisIt.

ORCID iDs
Xiaoshan Huang (黄小珊) https://orcid.org/0000-0003-2868-489X
Shane W. Davis https://orcid.org/0000-0001-7488-4468
Dong Zhang https://orcid.org/0000-0003-2582-4797

References
Aguirre, A., Hernquist, L., Katz, N., Gardner, J., & Weinberg, D. 2001, ApJL, 556, L11
Banda-Barragán, W. E., Parkin, E. R., Federrath, C., Crocker, R. M., & Bicknell, G. V. 2016, MNRAS, 455, 1309
Begelman, M. C., & Fabian, A. C. 1990, MNRAS, 244, 26P
Benson, A. J., Bower, R. G., Frenk, C. S., et al. 2003, ApJ, 599, 38
Bolatto, A. D., Warren, S. R., Leroy, A. K., et al. 2013, Nat, 499, 450
Bower, R. G., Benson, A. J., & Crain, R. A. 2012, MNRAS, 422, 2816
Brüggen, M., & Scannapieco, E. 2016, ApJ, 822, 31
Bruls, J. H. M. J., Vollmöller, P., & Schüssler, M. 1999, A&A, 348, 233
Chevalier, R. A., & Clegg, A. W. 1985, Nat, 317, 44
Cicone, C., Maiolino, R., Sturm, E., et al. 2014, A&A, 562, A21
Davis, S. W., Jiang, Y.-F., Stone, J. M., & Murray, N. 2014, ApJ, 796, 107
Draine, B. T. 2011, ApJ, 732, 100
Erb, D. K. 2008, ApJ, 674, 151
Fabian, A. C. 2012, ARA&A, 50, 455
Finlator, K., & Davé, R. 2008, MNRAS, 385, 2181
Geach, J. E., Tremonti, C., Diamond-Stanic, A. M., et al. 2018, ApJL, 864, L1
Heckman, T. M., & Best, P. N. 2014, ARA&A, 52, 589
Heckman, T. M., Lehnert, M. D., Strickland, D. K., & Armus, L. 2000, ApJS, 129, 493
Hopkins, P. F., Quataert, E., & Murray, N. 2012, MNRAS, 421, 3522
Ipavich, F. M. 1975, ApJ, 196, 107
Jiang, Y.-F., Stone, J. M., & Davis, S. W. 2014, ApJS, 213, 7
Kim, J.-G., Kim, W.-T., & Ostriker, E. C. 2016, ApJ, 819, 137
Kim, J.-G., Kim, W.-T., & Ostriker, E. C. 2018, ApJ, 859, 68
Klein, R. I., McKee, C. F., & Colella, P. 1994, ApJ, 420, 213
Krieger, N., Bolatto, A. D., Walter, F., et al. 2019, ApJ, 881, 43
Krumholz, M. R., & Thompson, T. A. 2012, ApJ, 760, 155
Krumholz, M. R., & Thompson, T. A. 2013, MNRAS, 434, 2329
Leroy, A. K., Walter, F., Martini, P., et al. 2015, ApJ, 814, 83
Martin, C. L. 1999, ApJ, 506, 222
Martin, C. L. 2005, ApJ, 621, 227
McCourt, M., O’Leary, R. M., Madigan, A.-M., & Quataert, E. 2015, MNRAS, 449, 2
Morganti, R., Veilleux, S., Oosterloo, T., Seng, S. H., & Rupke, D. 2016, A&A, 593, A30
Muñoz-Vergara, D., Nagar, N. M., Ramakrishnan, V., et al. 2019, MNRAS, 487, 3679
Murray, N., Ménard, B., & Thompson, T. A. 2011, ApJ, 735, 66
Murray, N., Quataert, E., & Thompson, T. A. 2005, ApJ, 618, 569
Namekata, D., Umemura, M., & Hasegawa, K. 2014, MNRAS, 443, 2018
Oosterloo, T., Raymond Oonk, J. B., Morganti, R., et al. 2017, A&A, 608, A38
Oppenheimer, B. D., Davé, R., Kereš, D., et al. 2010, MNRAS, 406, 2325
Pollack, J. B., Hollenbach, D., Beckwith, S., et al. 1994, ApJ, 421, 615
Proga, D., Jiang, Y.-F., Davis, S. W., Stone, J. M., & Smith, D. 2014, ApJ, 780, 51
Proga, D., & Waters, T. 2015, ApJ, 804, 137
Puchwein, E., & Springel, V. 2013, MNRAS, 428, 2966
Raskuti, S., Ostriker, E. C., & Skinner, M. A. 2016, ApJ, 829, 130
Rosen, A. L., Krumholz, M. R., McKee, C. F., & Klein, R. E. 2016, MNRAS, 463, 2553
Rupke, D. S., Veilleux, S., & Sanders, D. B. 2002, ApJ, 570, 588
Ruszkowski, M., Yang, H.-Y. K., & Zweibel, E. 2017, ApJ, 834, 208
Scannapieco, E., & Brüggen, M. 2015, ApJ, 805, 158
Schulz, R., Morganti, R., Nyland, K., et al. 2018, A&A, 617, A38
Semenov, D., Henning, T., Helling, C., Ilgner, M., & Sedlmayr, E. 2003, A&A, 410, 611
Shopbell, P. L., & Bland-Hawthorn, J. 1998, ApJ, 493, 129
Skinner, M. A., & Ostriker, E. C. 2013, ApJS, 206, 21
Skinner, M. A., & Ostriker, E. C. 2015, ApJ, 809, 187
Socrates, A., Davis, S. W., & Ramirez-Ruiz, E. 2008, ApJ, 687, 202
Spilker, J. S., Aravena, M., Béthermin, M., et al. 2018, Sci, 361, 1016
Strickland, D. K., & Heckman, T. M. 2007, ApJ, 658, 258
Strickland, D. K., & Heckman, T. M. 2009, ApJ, 697, 2030
Tadhunter, C., Holden, L., Ramos Almeida, C., & Batcheldor, D. 2019, MNRAS, 488, 1813
Thompson, T. A., Quataert, E., Zhang, D., & Weinberg, D. H. 2016, MNRAS, 455, 1830
Tsang, B. T.-H., & Milosavljević, M. 2015, MNRAS, 453, 1108
Veilleux, S., Cecil, G., & Bland-Hawthorn, J. 2005, ARA&A, 43, 769
Veilleux, S., Rupke, D. S. N., & Swaters, R. 2009, ApJL, 700, L149
Walter, F., Bolatto, A. D., Leroy, A. K., et al. 2017, ApJ, 835, 265
Waters, T., & Proga, D. 2016, MNRAS, 460, L79
Wiener, J., Zweibel, E. G., & Ruszkowski, M. 2019, MNRAS, 489, 205
Zhang, D. 2018, Galax, 6, 114
Zhang, D., & Davis, S. W. 2017, ApJ, 854, 110
Zhang, D., Thompson, T. A., Murray, N., & Quataert, E. 2014, ApJ, 784, 93