Modelling Star Formation and Stellar Feedback in Numerical Simulations of Galaxy Formation

A dissertation submitted for the degree of Doctor of Philosophy

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MATTHEW C. SMITH
MODELLING STAR FORMATION AND STELLAR FEEDBACK IN NUMERICAL SIMULATIONS OF GALAXY FORMATION

Summary

Remarkable progress has been made over the last few decades in furthering our understanding of the growth of cosmic structure. Nonetheless, there remains a great deal of uncertainty regarding the precise details of the complex baryonic physics that regulate galaxy formation. Any theory of star formation in galaxies must encompass the radiative cooling of gas into dark matter haloes, the formation of a turbulent, multiphase interstellar medium (ISM), the efficiency with which molecular gas is able to collapse into cores and ultimately stars, and the subsequent interaction of those stars with the gas through ionizing radiation, winds and supernova (SN) explosions. Given the highly non-linear nature of the problem, numerical simulations provide an invaluable tool with which to study galaxy formation. Yet, even with contemporary computational resources, the inherently large dynamical range of spatial scales that must be tackled makes the development of such models extremely challenging, inevitably leading to the adoption of ‘subgrid’ approximations at some scale. In this thesis, I explore new methods of incorporating the physics of star formation and stellar feedback into high resolution hydrodynamic simulations of galaxies.

I first describe a new implementation of star formation and SN feedback that I have developed for the state-of-the-art moving mesh code Arepo. I carry out a detailed study into various classes of subgrid SN feedback schemes commonly adopted in the literature, including injections of thermal and/or kinetic energy, two parametrizations of delayed cooling feedback and a ‘mechanical’ feedback scheme that injects the appropriate amount of momentum depending on the relevant scale of the SN remnant (SNR) resolved. All schemes make use of individually time-resolved SN events. Adopting isolated disk galaxy setups at different resolutions, with the highest resolution runs reasonably resolving the Sedov-Taylor phase of the SNR, I demonstrate that the mechanical scheme is the only physically well-posed method of those examined, is efficient at suppressing star formation, agrees well with the Kennicutt-Schmidt relation and leads to converged star formation rates and galaxy morphologies with increasing resolution without fine tuning any parameters. However, I find that it is difficult to produce outflows with high enough mass loading factors at all but the highest resolution. I discuss the various possible
solutions to this effect, including improved modelling of star formation.

Moving on to a more self-consistent setup, I carry out a suite of cosmological zoom-in simulations of low mass haloes at very high resolution, performed to $z = 4$, to investigate the ability of SN feedback models to produce realistic galaxies. The haloes are selected in a variety of environments, ranging from voids to crowded locations. In the majority of cases, SN feedback alone has little impact at early times even in low mass haloes ($\sim 10^{10} M_\odot$ at $z = 0$). This appears to be due largely to the build up of very dense gas prior to SN events, suggesting that other mechanisms (such as other stellar feedback processes) are required to regulate ISM properties before SNe occur. The effectiveness of the feedback also appears to be strongly dependent on the merger history of the halo.

Finally, I describe a new scheme to drive turbulence in isolated galaxy setups. The turbulent structure of the ISM very likely regulates star formation efficiencies on small scales, as well as affecting the clustering of SNe. The large range of potential drivers of ISM turbulence are not fully understood and are, in any case, unlikely to arise \textit{ab initio} in a whole galaxy simulation. I therefore neglect these details and adopt a highly idealised approach, artificially driving turbulence to produce an ISM structure of my choice. This enables me to study the effects of a given level of ISM turbulence on global galaxy properties, such as the fragmentation scale of the disk and the impact on SN feedback efficiencies. I demonstrate this technique in the context of simulations of isolated dwarfs, finding that moderate levels of turbulent driving in combination with SN feedback can produce a steady-state of star formation rates and global galaxy properties, rather than the extremely violent SN feedback that is produced by a rapidly fragmenting disk.
DECLARATION OF ORIGINALITY

I, Matthew Carey Smith, declare that this thesis entitled ‘Modelling star formation and stellar feedback in numerical simulations of galaxy formation’, and the work presented in it, is the result of my own research and includes nothing which is the outcome of work done in collaboration except as declared in the Preface and specified in the text.

Chapter 2 is based on work completed in collaboration with D. Sijacki and S. Shen, published as: Smith, M. C., Sijacki, D., Shen, S., 2018, ‘Supernova feedback in numerical simulations of galaxy formation: separating physics from numerics’, MNRAS, 478, 302

I confirm that this thesis is not substantially the same as any that I have submitted, or, is being concurrently submitted for a degree or diploma or other qualification at the University of Cambridge or any other University or similar institution except as declared in the Preface and specified in the text. I further state that no substantial part of my dissertation has already been submitted, or, is being concurrently submitted for any such degree, diploma or other qualification at the University of Cambridge or any other University or similar institution. The length of this thesis does not exceed the stated limit of the Degree Committee of Physics and Chemistry of 60,000 words.

Matthew C. Smith
Cambridge, April 2018
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To my grandfathers.
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1. Introduction

1.1 Cosmic structure formation

Large strides in observational cosmology over the last few decades have transformed our understanding of the observable Universe, resulting in the development of a standard model of cosmic structure formation. Einstein’s theory of general relativity provides the framework with which we can describe the dynamics of the universe. Observational evidence shows that the universe obeys the cosmological principle (i.e. it is homogeneous and isotropic) on large scales, allowing it to be described by the Robertson-Walker metric, leading to the derivation of the Friedmann equations to describe its evolution (this is described in more detail in the next section).

We know that we live in an expanding universe that originated from an initially hot, dense state. The cosmic microwave background (CMB), produced as photons decoupled from matter approximately $4 \times 10^5$ yr after the Big Bang and subsequently redshifted into the microwave range by the expansion of the universe, contains an imprint of the structure of the universe at those early times. By examining temperature fluctuations in the CMB, constraints can be placed on the cosmological parameters that describe the evolution of the universe. Fig. 1.1 shows the power spectrum of the temperature fluctuations of the CMB as measured by the Planck satellite (Planck Collaboration XIII 2016), while Table 1.1 shows the derived cosmological parameters. Measurements of the CMB, galaxy clustering and Type Ia SNe have enabled the development of a well constrained model of cosmology, known as $\Lambda$CDM. The total critical density (defined later) is measured to be close to unity, meaning the universe has a flat geometry. The expansion of the universe is accelerating, driven by an energy density component of the universe that behaves as if it has negative pressure, making up $\sim 70\%$ of the total energy density. Determining the exact nature of this ‘dark energy’ is currently a major challenge facing modern cosmology. This component appears in the Friedmann equations as a non-zero cosmological constant $\Lambda$.

Matter makes up $\sim 30\%$ of the energy density, but only $\sim 16\%$ of this is in the form of baryonic matter. The remaining fraction, ‘dark matter’, is extremely weakly interacting (effectively collisionless) and is currently believed to have had
negligible thermal velocities when it decoupled. The resulting free-streaming length impacts the degree to which the collisionless dark matter would have washed out small scale structure. In a universe with cold dark matter (CDM), the presently favoured model, structures form in a ‘bottom-up’ fashion, with initially small density fluctuations (of order $10^{-5}$ at recombination) growing via gravitational instability. Structure then forms in an hierarchical fashion, with small dark matter haloes that formed first merging to form larger structures. The primary driver of structure formation is therefore gravity. However, while the collisionless dark matter may make up the dominant component of the matter in the universe, to completely understand the formation and evolution of cosmic structures, we must also examine the multitude of complex processes that affect baryons.

1.1.1 The homogeneous universe

Given the difficulty involved in solving Einstein’s field equations (which relate the geometrical properties of spacetime to its stress-energy tensor) for arbitrary matter distributions and the general lack of empirical data about the large scale structure of the universe, the obvious starting point for early cosmologists was to construct very simple models. The assumptions that the universe was homogeneous and
Table 1.1 Base ΛCDM parameters from Planck Collaboration XIII (2016) with 68% confidence limits for the fit to the CMB power spectrum in combination with lensing reconstruction. From left to right, the parameters are the matter density parameter, the physical baryon matter density parameter, the dark energy density parameter, the fluctuation amplitude at $8h^{-1}\text{Mpc}$, the present day Hubble parameter and the age of the universe. Note that some of these parameters are derived from the actual 6 independent parameters used in the fit; these represent the most intuitive choice of parameters to describe the cosmology.

<table>
<thead>
<tr>
<th>$\Omega_\text{m}$</th>
<th>$\Omega_\text{b}h^2$</th>
<th>$\Omega_\Lambda$</th>
<th>$\sigma_8$</th>
<th>$H_0/(\text{km s}^{-1}\text{Mpc}^{-1})$</th>
<th>$t_0/\text{Gyr}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.308</td>
<td>0.02226</td>
<td>0.692</td>
<td>0.8149</td>
<td>67.81</td>
<td>13.799</td>
</tr>
<tr>
<td>$\pm 0.012$</td>
<td>$\pm 0.00023$</td>
<td>$\pm 0.012$</td>
<td>$\pm 0.0093$</td>
<td>$\pm 0.92$</td>
<td>$\pm 0.038$</td>
</tr>
</tbody>
</table>

isotropic, together known as the cosmological principle, allow such investigations to be tractable.

Before proceeding further, we must ask to what extent does the universe obey the cosmological principle? Obviously, on the scale of individual galaxies the universe is neither homogeneous nor isotropic. Nonetheless, the large-scale ($\sim 100\text{ Mpc}$) distribution of galaxies and galaxy clusters seems to obey these properties, as does the distribution of quasars. Finally, measurements of temperature fluctuations in the CMB show that the universe is anisotropic on large scales by no more than about 1 part in $10^5$.

The most general spacetime metric describing a universe in which the cosmological principle holds is the Robertson-Walker metric:

$$ds^2 = (c dt)^2 - a(t)^2 \left[ \frac{dr^2}{1 - Kr} + r^2 \left( d\theta^2 + \sin^2 \theta d\phi^2 \right) \right],$$  \hspace{1cm} (1.1)

where $r$, $\theta$ and $\phi$ are the comoving coordinates and $t$ is the proper time. The expansion factor $a(t)$ encodes the growth of the spatial extent of the universe and is normally defined such that $a(t_0) = 1$ at the present time. The constant $K$ is known as the curvature parameter; it can be scaled in such a way that it takes on one of three values. $K = -1$, 0 and 1 produces an open (hyperbolic), flat (Euclidean) or closed (spherical) geometry, respectively. Under the assumption that the universe can be described by a homogeneous and ideal fluid, we can solve the field equations with the Robertson-Walker metric to produce the Friedmann equations:
\[ \frac{\dot{a}}{a} = -\frac{4\pi G}{3} \left( \rho + \frac{3p}{c^2} \right) + \frac{\Lambda c^2}{3}, \] (1.2)

\[ \left( \frac{\dot{a}}{a} \right)^2 = \frac{8\pi G}{3} \rho - \frac{K c^2}{a^2} + \frac{\Lambda c^2}{3}, \] (1.3)

where we have allowed for a non-vanishing cosmological constant \( \Lambda \). It can be seen that the evolution of the expansion parameter can be obtained by solving the Friedmann equations if we have an equation of state relating the pressure \( p \) to the density \( \rho \). Typically, the equation of state is parametrised as

\[ p = w \rho c^2, \] (1.4)

where \( w \) is in the interval \([0, 1]\). In the early stages of its evolution, the universe was radiation dominated; \( w = 1/3 \) corresponds to a fluid of relativistic, non-degenerate particles in thermodynamic equilibrium. On the other hand, \( w = 0 \) is a good approximation for an ideal, non-relativistic fluid, relevant for a matter dominated universe.

A convenient derived parameter is the Hubble parameter, \( H = \dot{a}/a \). The recession velocity of some source relative to observer due to the background expansion of the universe is \( v_r = (\dot{a}/a)a = Hr \). The currently estimated value of the Hubble parameter, \( H_0 \), is \( 70 \text{ km s}^{-1} \text{ Mpc}^{-1} \). It is often useful to define a dimensionless factor \( h = H_0/100 \text{ km s}^{-1} \text{ Mpc}^{-1} \) to absorb uncertainties in the measurement of \( H_0 \).

The critical density of the universe, \( \rho_{\text{crit}} \), corresponds to the density that would produce a flat universe. We can express the present day value as

\[ \rho_{\text{crit},0} = \frac{3H_0^2}{8\pi G}. \] (1.5)

It is convenient to define the total density parameter at the present time as

\[ \Omega_{\text{tot},0} = \frac{\rho_0}{\rho_{\text{crit},0}}, \] (1.6)

such that \( \Omega_{\text{tot},0} > 1, = 1 \) or \( < 1 \) corresponds to closed, flat and open geometries, respectively. From the Friedmann equations and the adiabatic expansion of the universe, the relation between the density of a component having a particular value of \( w \) and the expansion factor is

\[ \rho_w a^{3(1+w)} = \rho_{w,0} a_0^{3(1+w)}. \] (1.7)
Then, expressing the contribution to the total density parameter for a given component as $\Omega_{w,0} = \rho_0 / \rho_{w,0}$, the time evolution of the Hubble parameter can be seen to be

$$H^2(t) \equiv \left(\frac{\dot{a}}{a}\right)^2 = H_0^2 \left(\frac{a_0}{a}\right)^2 \left\{ \Omega_{w,0} \left(\left(\frac{a_0}{a}\right)^{1+3w} - 1\right) \right\} + 1 ,$$

(1.8)

where the sum is over all components. For a flat universe in which only matter and the cosmological constant are significant contributors to the total cosmic energy density (as in $\Lambda$CDM), this simplifies to

$$H^2(t) = H_0^2 \left[ \Omega_{m,0} \left(\frac{a_0}{a}\right)^3 + \Omega_{\Lambda,0} \right] ,$$

(1.9)

where the cosmological constant has been assumed to act as a fluid with a $w = -1$ equation of state. It is useful at this stage to define the variable $z = \frac{a_0}{a} - 1$, known as the redshift. This essentially expresses the ratio between the present day expansion factor and its value at some earlier epoch. It is convenient from an observational perspective because it can be measured from the spectra of an object as

$$z = \frac{\lambda_{\text{obs}} - \lambda_{\text{em}}}{\lambda_{\text{em}}} ,$$

(1.10)

where $\lambda_{\text{obs}}$ is the observed wavelength of light that was emitted in the rest frame of the object at a wavelength of $\lambda_{\text{em}}$ at time $t$. Having corrected for any peculiar (i.e. inherent) motions of the object, this therefore provides a distance measurement and correspondingly a lookback time.

### 1.1.2 Linear theory of structure growth

The equations governing an ideal, non-relativistic fluid with gravity are:

1. The mass conservation equation,

$$\frac{\partial \rho}{\partial t} + \nabla (\rho \mathbf{v}) = 0 .$$

(1.11)

2. The Euler equation,

$$\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla) \mathbf{v} + \frac{1}{\rho} \nabla p + \nabla \Phi = 0 .$$

(1.12)

3. The equation of state,

$$p = p(\rho, S) .$$

(1.13)
4. The Poisson equation,
\[ \nabla^2 \Phi = 4\pi G \rho. \]  
(1.14)

5. The entropy evolution,
\[ \frac{dS}{dt} = 0. \]  
(1.15)

A trivial solution to these equations is that of an homogeneous expanding (or contracting) fluid:
\[ \rho = \rho_0 \left( \frac{a}{a_0} \right)^{-3}, \]  
(1.16)
\[ v = -\frac{\dot{a}}{a} r = H r, \]  
(1.17)
\[ \Phi = \frac{2\pi G}{3} \rho r^2, \]  
(1.18)
where \( \rho_0 \) is the density of the fluid at some time with expansion factor \( a_0 \). However, while the cosmological principle tells us that the universe is homogeneous on large scales, it contains primordial density fluctuations that will seed future structure. Fortunately, these initial fluctuations are very small, \( \delta \rho/\rho \sim 10^{-5} \), so they can initially be treated as perturbations to a homogeneous background, allowing us to approximate their evolution with linear perturbation theory. This approximation is valid as long as the density contrast relative to the background density, \( \rho_b \), remains significantly less than unity:
\[ \delta = \frac{\rho - \rho_b}{\rho_b} \ll 1. \]  
(1.19)

We can now express the density, velocity, pressure, potential and specific entropy in terms of a uniform background component and a small perturbation:
\[ \rho = \rho_b + \delta \rho \]  
(1.20)
\[ \mathbf{v} = H \mathbf{r} + \mathbf{u} \]  
(1.21)
\[ p = p_b + \delta p \]  
(1.22)
\[ \Phi = \Phi_b + \delta \Phi \]  
(1.23)
\[ s = s_b + \delta s. \]  
(1.24)

Inserting these into the fluid equations and subtracting the homogeneous equations we obtain the perturbation equations:
\[
\frac{\partial (\delta \rho)}{\partial t} + \frac{3 \dot{a}}{a} \delta \rho + \frac{\dot{a}}{a} (\mathbf{r} \cdot \nabla) \delta \rho + \rho_b (\nabla \cdot \mathbf{u}) = 0
\]  
(1.25)

\[
\frac{\partial \mathbf{u}}{\partial t} + \frac{\dot{a}}{a} (\mathbf{r} \cdot \nabla) \mathbf{u} + \frac{1}{\rho_b} \nabla \delta p + \nabla \delta \Phi = 0
\]  
(1.26)

\[
\nabla^2 \delta \Phi = 4 \pi G \delta \rho
\]  
(1.27)

\[
\frac{\partial (\delta s)}{\partial t} + \frac{\dot{a}}{a} (\mathbf{r} \cdot \nabla) \delta s = 0.
\]  
(1.28)

Transforming to Fourier space and performing some algebra, one obtains the following equation for the density contrast:

\[
\ddot{\delta} + 2 \frac{\dot{a}}{a} \dot{\delta} + \left( c_s^2 k^2 - 4 \pi G \rho_b \right) \delta = 0,
\]  
(1.29)

where \(k\) is the (proper) wavenumber and \(c_s\) is the sound speed, which is obtained as the square root of the isentropic derivative of the pressure with respect to density. Considering plane wave solutions, \(\delta \propto \exp (i \omega t)\), a non-zero solution is only possible if the following dispersion relation is satisfied:

\[
\omega^2 - c_s^2 k^2 + 4 \pi G \rho_b = 0.
\]  
(1.30)

There are two classes of solutions, depending on whether the wavelength \(\lambda = 2\pi/k\) is greater or less than

\[
\lambda_J = c_s \left( \frac{\pi}{G \rho_b} \right)^{1/2},
\]  
(1.31)

which is known as the Jeans length.\(^1\) If \(\lambda < \lambda_J\), then \(\omega\) is real, producing solutions which correspond to sound waves. If, however, \(\lambda > \lambda_J\), then \(\omega\) is imaginary, representing a stationary wave of either increasing or decreasing amplitude. This is the gravitational or Jeans instability. For \(\lambda \gg \lambda_J\), the characteristic timescale for the evolution of the amplitude of the perturbation coincides with the gravitational free-fall time, \(\tau_{ff} \sim 1/\sqrt{G \rho_b}\). A physically intuitive way of understanding the Jeans length is to consider it as representing the scale on which the gravitational force begins to dominate the pressure force. That is to say that the gravitational free-fall time is less than the sound crossing time, \(\tau_h \sim \lambda/v_s\). While we have examined the Jeans length here in the context of cosmological density fluctuations, it is also a useful concept when examining gravitational collapse on all scales, for example in star forming regions.

\(^1\)An analogous analysis also holds in the case of a completely collisionless fluid, in which case (for a Maxwellian distribution of velocities) the sound speed is replaced by the velocity dispersion of the fluid, \(\sigma\).
1.1.3 A non-linear example: the spherical ‘top-hat’ collapse

As described above, once the magnitude of a density perturbation becomes comparable to the background density (i.e. when $\delta \ll 1$ is no longer satisfied), we can no longer make use of perturbation theory. However, it is still possible to follow the evolution of structures into the strongly non-linear regime analytically in a few simple cases. For example, let us consider a spherical top-hat perturbation (i.e. the density is constant within the perturbation) with an initial perturbation $\delta_i > 0$ and $|\delta_i| \ll 1$ (such that it initially satisfies linear theory) at time $t_i$ with negligible pressure and no peculiar velocity. This particular situation allows us to treat the perturbation as a separate universe. For simplicity, we can assume that the underlying universe at $t_i$ is described by a flat, matter dominated (i.e. $\Omega_m = 1$) universe; an Einstein-de Sitter model. In such a universe, the solution to eq. 1.29 is composed of two modes: a growing mode, $\delta_+ \propto t^{2/3} \propto a$, and a decaying mode, $\delta_- \propto t^{-1} \propto a^{-3/2}$. This means that we can express our density perturbation as

$$\delta = \delta_+ (t_i) \left( \frac{t}{t_i} \right)^{2/3} + \delta_- (t_i) \left( \frac{t}{t_i} \right)^{-1} .$$  \hspace{1cm} (1.32)

It can be seen that after some time $\delta_+ (t_i) \left( t/t_i \right)^{2/3} \gg \delta_- (t_i) \left( t/t_i \right)^{-1}$, while enforcing our boundary condition that the peculiar velocity is zero leads to $\delta_+ (t_i) = \frac{3}{5} \delta_i$. Our ‘sub-universe’ inside the perturbation evolves like a Friedmann model with an initial density parameter $\Omega_p = \Omega (t_i) (1 + \delta_i)$ where $\Omega (t_i)$ is the density parameter of the unperturbed universe. For the perturbation to decouple from the general expansion of the universe and collapse, we require $\Omega_p > 1$. This leads to the constraint on the growing mode that

$$\delta_+ (t_i) = \frac{3}{5} \delta_i > \frac{3}{5} \frac{1 - \Omega (t_i)}{\Omega (t_i)} = \frac{3}{5} \frac{1 - \Omega_0}{\Omega_0 (1 + z_i)},$$  \hspace{1cm} (1.33)

where on the right hand side we have substituted in the redshift evolution of the density parameter in terms of the present day value. Therefore, it can be seen that for $\Omega_0 \geq 1$ every perturbation will collapse. On the other hand, if $\Omega_0 < 1$ then $\delta_i$ must be large enough such the perturbation enters the non-linear regime before the universe becomes curvature dominated and begins undecelerated free expansion.

To compare the loss in accuracy experienced by applying linear theory in this regime, we can examine the values for the density of the perturbation at maximum compression calculated in the framework given above versus an extrapolation of
linear theory. The former predicts the ratio of the density of the perturbation relative to the background universe of $\sim 180$, whereas extrapolating linear theory produces a value of $\sim 1.68$. Thus, it is apparent that while linear theory can help us understand the early stages of structure formation, non-linear theories are required to help us understand the structures that will eventually host galaxies. Unfortunately, while the exact solutions arising from simple models such as that presented in this section can give insights into non-linear structure formation, to tackle the growth of real structure we must make use of numerical methods.

1.2 Numerical methods

As we described in the previous section, accurate analytic calculations in the non-linear regime rapidly become intractable for all but the simplest configurations. Therefore, with the advent of increasingly powerful computational resources, numerical methods have become an important tool for studying the growth of cosmic structure and galaxy formation over the last few decades. While a great deal of progress has been made simulating structure formation under the influence of gravity alone, the focus of modern techniques is on coupling this with the complex range of baryonic processes that affect galaxy evolution. The eventual aim of research in this area is to accurately model all relevant physical processes from first principles (i.e. ab initio) in a self-consistent manner. Two major obstacles to achieving this goal are typically a lack of dynamic range (impacting the spatial resolution that can be achieved) and an incomplete understanding of key astrophysical processes (for example, the small scale details of supermassive blackhole accretion, mergers and feedback). Both of these issues are usually dealt with by adopting ‘subgrid’ schemes that include the effects of unresolved or missing processes. These may be designed to replicate results from analytic or higher resolution numerical models. Alternatively, they may be tuned to match observations in a phenomenological manner. In addition, computational resources may limit the full range of physics that can be included in a numerical simulation. For example, we now know that the effects of radiation and magnetic fields are likely to be important in a wide range of scenarios, but the cost of solving the full set of radiation and/or magneto- hydrodynamic equations can potentially be prohibitively high so they may be omitted or approximated. Nonetheless, with a combination of the increasing performance of computer hardware and the continued development of more sophisticated algorithms, remarkable progress has continued to be made in numerical simulations of galaxy formation over the last few years.
1.2.1 N-body techniques

An obvious way to model the time evolution of a collisionless self-gravitating three dimensional fluid is to discretize it into particles and then treat the system as an N-body problem. As long as the mass resolution is much smaller than the mass-scale of interest in the system being studied, this is a reasonable approximation. The simplest numerical method to solve an N-body problem is that of direct summation, often referred to as a particle-particle (PP) method: one obtains the total gravitational force acting on a given particle as the sum of the forces from every other individual particle. However, the approximation of the collisionless fluid as an ensemble of particles is valid only in the absence of two-body interactions (such as large angle scattering) between particles that would not occur in an extended matter distribution. For this reason, the gravitational force must be ‘softened’ on some length scale, $\epsilon$. A simple method of achieving this is to modify the force law between two particles as

$$| \mathbf{F} | = \frac{Gm_1m_2}{r^2} \rightarrow \frac{Gm_1m_2}{(r^2 + \epsilon^2)},$$

(1.34)

although more sophisticated forms are often used. Empirically, typical values of between $1/50 - 1/25$ of the mean particle separation are appropriate.

While the PP method may be conceptually and algorithmically simple, it suffers from the major drawback of being somewhat computationally expensive: for obvious reasons, the time taken to integrate the forces between every pair of particles is proportional to $N^2$. For particle numbers $\gtrsim 10^6$, this becomes intractable even with contemporary computing resources. One alternative to the PP method is the particle-mesh (PM) method (e.g. Klypin & Shandarin 1983). The key point here is that the Poisson equation is easier to solve in Fourier space:

$$\nabla^2 \Phi = 4\pi G \rho \rightarrow \hat{\Phi} = -4\pi G \frac{\hat{\rho}}{k^2},$$

(1.35)

where $\hat{\Phi}$ and $\hat{\rho}$ are the Fourier transformed potential and density, respectively. The density field is discretized from the particles onto a regular grid, a Fast Fourier Transform (FFT) is applied, the potential is calculated in Fourier space and then transformed back onto a regular grid with periodic boundary conditions. This produces more favourable scaling with particle number than the PP method, but the force resolution is limited to the size of the grid. A hybrid approach, the P$^3$M method (Efstathiou et al. 1985; Couchman 1991) balances computational time and accuracy by using the PP to calculate the short range forces between individual
Fig. 1.2 The Millennium-XXL is the largest cosmological N-body simulation ever performed. *Left:* Density slices focusing on the most massive halo in the simulation at $z = 0$, starting with the full 4.1 Gpc box and zooming to an 8.1 Mpc region. The structure of the cosmic web is similar across many orders of magnitude. *Right:* The halo mass function for the simulation, along with those from the predecessor Millennium and Millennium-II simulations. Together, they cover 8 decades in mass.

Particles while using the PM method for long range forces.

An alternative approach is taken by tree algorithms (Barnes & Hut 1986), in which particles are grouped together according to their distance from the particle for which the force is being determined. The force from the groups can then be replaced by its multipole expansion. This is typically achieved by repeatedly subdividing the computational domain into a sequence of nested cubes which function as nodes for an oct-tree (in 3 dimensions) structure. The force acting on a particle can then be evaluated by walking the tree and integrating the contributions from the various nodes. On the smallest scale, this will be individual PP forces while with increasing distance they will be summed contributions of the particles within a node. One can specify a maximum opening angle, $\theta$, to control the level of accuracy used such that a node will be ‘opened’ and its eight subnodes treated individually in the calculation if $r > l/\theta$ where $r$ and $l$ are the distance to and diameter of the node, respectively. This algorithm theoretically allows the force computation to scale as $O(N \log N)$ rather than the $O(N^2)$ of the PP method. While the tree method is able to adapt to the particular clustering state of a given simulated system, the major sources of computational expense occur if the tree becomes too deep due to a large density contrast (meaning tree walks can take too many steps) or if large fractions of the tree must be reconstructed regularly.
These issues can be mitigated to some extent by using a hybrid TreePM scheme where the longest range force are calculated using a PM method (see e.g. Xu 1995; Bode et al. 2000; Bagla 2002; Springel 2005). Fig. 1.2 shows an example of a large scale cosmological N-body simulation making use of a TreePM scheme. The Millennium-XXL simulation (Angulo et al. 2012), along with its predecessors, were run to enable the investigation of the impact of ΛCDM cosmology on the resulting structure formation (which cannot be followed with non-linear techniques) and to provide predictions for comparison to observations, such as the spatial distribution of galaxies.

1.2.2 Numerical hydrodynamics

Because dark matter is the dominant component of mass in the universe, modelling the growth of cosmic structure as a collisionless fluid as above is a reasonable approximation for large scales. However, if one is concerned with smaller scale evolution, particularly in the study of galaxy formation, a treatment of numerical hydrodynamics is required in addition to the computation of self-gravity described above in order to model baryonic components of matter. The task is essentially to solve the fluid equations (eq. 1.11 - 1.15).

Eulerian codes

In an Eulerian code, the computational domain is divided into a Cartesian\(^2\) grid. The hydrodynamic equations are solved in their Eulerian form by computing the fluxes between neighbouring cells. While using a grid of a fixed resolution is the simplest approach, by nesting higher resolution grids in areas of higher density, usually ‘on-the-fly’, the dynamic range that can be treated accurately by the code can be substantially increased; this technique is known as adaptive mesh refinement (AMR). While Eulerian codes do not suffer from some of the drawbacks of traditional Lagrangian methods (described below), they can potentially have a few disadvantages. They are not Galilean invariant (meaning that they are sensitive to bulk velocities), they do not guarantee conservation of angular momentum and they may experience advection errors arising from numerical diffusion. However, these issues are resolution dependent and can therefore be largely mitigated by efficient, contemporary AMR codes.

\(^2\)Although other coordinate systems may be adopted depending on underlying symmetries in the problem at hand.
Lagrangian codes

In a Lagrangian code, rather than discretizing the computational volume, the fluid itself is discretized, being represented by a set of particles. These act as interpolants of the continuum fluid flow. The most common technique used in this type of approach is known as smoothed-particle hydrodynamics (SPH) (Lucy 1977; Gingold & Monaghan 1977). Each SPH particle is assigned a position, velocity, (fixed) mass and other relevant internal thermodynamic properties. The value of any field of the fluid, $A$, at an arbitrary position $r$ is given as a mass and kernel-weighted sum over the particles:

$$A(r) = \sum_{j=1}^{N} \frac{m_j A_j}{\rho_j} W(r - r_j, h),$$  \hspace{1cm} (1.36)

where $W(r, h)$ is a smoothing kernel of radius $h$. The form of the kernel is usually chosen to have finite support (for example the cubic spline) such that in practice, the weighted sum is over particles within a radius $\sim h$. The spatial derivative of the smoothed field is then

$$\nabla A(r) = \sum_{j=1}^{N} \frac{m_j A_j}{\rho_j} \nabla W(r - r_j, h).$$  \hspace{1cm} (1.37)

Eq. 1.36 and 1.37 can then be used to solve the fluid equations for the changes in density, velocity and internal energy of a particle.

While fixed values of $h$ can be adopted, one of SPH’s greatest strengths is revealed when $h$ is allowed to vary to enclose some fixed number of neighbours. The result is that the spatial resolution of the code changes continuously with density, meaning that it can operate efficiently over a large dynamic range. Also, because of its Lagrangian nature, SPH codes guarantee exact conservation of mass, energy, entropy, linear and angular momentum. The use of particles for the discretization of the fluid also means that SPH couples to existing N-body gravity solvers in a natural way. Traditional SPH methods provide poor treatment of shocks and are typically unable to resolve hydrodynamic instabilities such as Kelvin-Helmholtz and Rayleigh-Taylor. However, modern schemes are able to avoid these adverse effects (see e.g. Price 2008; Hopkins et al. 2013; Hobbs et al. 2013).

A Hybrid Eulerian-Lagrangian approach: Arepo

Given that Eulerian and Lagrangian codes each have distinct advantages, various hybrid schemes have been proposed. An obvious direction in which to proceed is to attempt to solve Eulerian hydrodynamics on a moving mesh, for example by
allowing a Cartesian grid to deform (see e.g. Gnedin 1995; Pen 1998). However, the mesh will tend to become ‘tangled’ as cell shapes become heavily distorted. One method to deal with this is to remap the distorted mesh onto a more regular mesh to avoid this issue, a feature of arbitrary Lagrange-Eulerian (ALE) approaches. Unfortunately, remapping is unavoidably a diffusive process and the imposing of regular meshes does not help with the issue of the lack of Galilean invariance common to Eulerian approaches.

The code AREPO (Springel 2010) sidesteps these issues by solving the hydrodynamics on an unstructured moving mesh defined as the Voronoi tessellation of a set of discrete points. The mesh generating points are analogous to SPH particles. If they are allowed to drift with the local fluid velocity, then the resulting mesh moves smoothly with the flow (the use of a Voronoi tessellation avoids issues of mesh tangling). This means that the code is Galilean invariant and adjusts its spatial resolution automatically and continuously. The scheme also inherits the inherent high accuracy of Eulerian methods in treating shocks and contact discontinuities and has been demonstrated to resolve fluid instabilities. However, it is still not guaranteed to conserve angular momentum and, like any grid based scheme, can potentially experience numerical diffusion to some extent.

If an arbitrary distribution of mesh generating points is used, the resulting mesh will be unstructured but also irregular; the latter point means that cells may be distorted and suffer from large aspect ratios, which is an undesirable consequence. Lloyd (1982) gives a simple algorithm to obtain a Voronoi tessellation with regular shaped cells. Iteratively moving the mesh generating points to the centres of mass of the existing cells then constructing a new mesh results in regular, rounded cell geometries. By adding an additional drift velocity (proportional to the fluid sound speed) to the mesh generating points towards the cell centres of mass in addition to their normal motion with the fluid, a regular mesh geometry is maintained.

The combination of the motion of the mesh generating points with the fluid and the cell regularisation techniques means that cells tend to keep a relatively constant mass. However, some scatter in cell masses will inevitably occur after some time, particularly in complicated flows and shocks. To this end, AREPO can refine or derefine cells by adding or removing mesh generating points. In the former case, a single point is split into two which are initially positioned close to each other; the mesh regularisation scheme will move them apart over the next few timesteps. In the case of derefinement, a mesh generating point is simply removed and its various quantities (mass, momentum energy etc.) distributed to its neighbours. Upon the next mesh construction, the ‘hole’ where the cell used to be is automatically filled.
in. The standard refinement strategy in AREPO is to keep cell masses within a factor of 2 of some target mass, but any arbitrary criteria can be applied. For example, one can refine to resolve the local Jeans length (e.g. Simpson et al. 2016) or to increase spatial resolution around an object of interest such as black holes (Curtis & Sijacki 2015).

1.2.3 Additional physics

While the methods described above allow the modelling of a self-gravitating, ideal fluid, they cannot capture the broad range of astrophysical phenomena which are also necessary for studying galaxy formation. Additional physics must be added on top of the pure gravity and hydrodynamics solvers, with varying degrees of approximation.

So far, we have only considered a non-radiative fluid. In reality, gas in astrophysical scenarios is subject to a number of cooling and heating processes. Once gas is allowed to radiatively cool, it will be able to flow into the centre of the potential wells of dark matter haloes. Atomic gas composed of hydrogen and helium cools due to free-free emission, collisional excitation and ionization, recombination and Compton cooling from CMB photons (at high redshifts). The strength of the various contributions from each process are determined by the temperature and the abundances of the different ions. If the gas is in ionization equilibrium, then cooling rates normalised by the square of the density can be tabulated as a function of temperature. Otherwise, the time dependent ionization equations must be solved to calculate the state of the gas at any given moment. If the timescale for the radiative processes is short compared to the hydrodynamic timescale, then the equilibrium values can be assumed (this is usually the case for atomic gas), which means that the cooling rates as a function of temperature can be incorporated into the fluid equations. In addition to primordial species, cooling from metal lines is important in enriched gas. Fig. 1.3 shows cooling functions as a function of temperature for gas of varying levels of metallicity (Sutherland & Dopita 1993). Molecular gas, of which molecular hydrogen (H₂) is the simplest, possesses vibrational and rotational states in addition to electronic levels. The energy involved in one of these transitions is much lower than those atomic transitions, so molecules can be rotationally or vibrationally excited even in low temperature gas. H₂ therefore acts as a coolant in very cold, dense gas. Once again, the various abundances of relevant species must be calculated from their formation and destruction rates (which will depend on density and temperature). This typically must proceed with a full non-equilibrium calculation.
Introduction

Fig. 1.3 Cooling functions above $10^4 \text{K}$ as a function of temperature for gas of varying levels of metallicity from Sutherland & Dopita (1993). The shape of the metal-free cooling function is dominated by peaks at $1.5 \times 10^4 \text{K}$ and $10^5 \text{K}$ due to collisionally excited levels of hydrogen and helium, respectively, before Bremsstrahlung (free-free) emission creates a rising tail at high temperatures. In an enriched gas, metals such as oxygen, carbon and nitrogen contribute to a peak at $10^5 \text{K}$ while metals such as neon, iron and silicon enhance the cooling rates at $\sim 10^6 \text{K}$. Below $10^4 \text{K}$ (not shown) cooling is dominated by metals and molecular species.

In addition to cooling processes, there are a variety of heating processes that affect gas. For example, gas may be photoionized either by local stars or by the general UV background produced from star-forming galaxies and/or quasars that pervades the universe during and after cosmic reionization ($z \lesssim 6$). Photoionization has two effects: it eliminates line excitation and ionization as cooling processes (recombination cooling increases, but the net effect is to reduce the total cooling rate). Secondly, it heats gas as photoelectrons carry off residual energy. In the case of molecular hydrogen, photodissociation destroys the $\text{H}_2$ molecules and removes them as a coolant. Dense gas may be able to resist the effects of photoionizing photons through self-shielding. The interaction of radiation with gas is a complex process and requires some level of approximation ranging from potentially expensive radiation hydrodynamic (RHD) solvers to simpler subgrid models. The UV background, for instance, is usually treated as a time-dependent but spatially
uniform heating and ionizing term in pure hydrodynamic simulations.

As well as photoionizing gas, massive stars can also provide mechanical input into the surrounding medium through stellar winds and radiation pressure before finally exploding as SNe. In addition, the shock fronts of expanding supernova remnants (SNRs) are believed to be a major source of cosmic rays, which themselves affect gas by providing a heat source and additional pressure. Of course, the basic gravity and hydrodynamic solvers we have discussed above cannot provide a treatment of these phenomena completely a priori, so this missing physics must be included at some level of approximation. While we leave a full discussion of how these sources of stellar feedback can be modelled until Chapter 2, it is worth noting here that this is a frequent source of uncertainty in galaxy formation simulations. The same applies to techniques used to model star formation itself in simulations; as the stellar masses formed in galaxies are both an observable quantity and a source of feedback, it is very important to use an accurate treatment of the relevant processes (we leave a detailed discussion about numerical techniques in this area until subsequent chapters). The final major source of feedback comes from active galactic nuclei (AGN), which are believed to regulate massive galaxies and galaxy properties by releasing large quantities of energy. AGN are not the focus of this work, but they are another primary example of additional physics that must be included in numerical hydrodynamic simulations if they are to accurately describe galaxy formation.

As we have briefly demonstrated in the previous paragraphs, there exists a large range of additional astrophysical phenomena that must be included ‘by hand’ to some degree. Such models are often referred to as ‘subgrid physics’. The level of approximation (or degree of abstraction from the underlying physics) that must be adopted is often dependent on the scale to be modelled. For example, in a large cosmological simulation there is no prospect of resolving the ISM and relevant processes such as star formation and stellar feedback, so the results of these processes must be added in a phenomenological manner tuned to observations (for example, driving galactic winds at a certain rate relative to halo mass). In simulations with higher resolution, processes can be modelled more explicitly because the pure hydrodynamic solver can model the physics in a natural manner. Some forms of missing physics, on the other hand, are likely to always remain an abstraction in galaxy formation simulations; for example, as it is unlikely that codes will ever contain the pertinent physics to model stellar evolution from first principles, we must use approximations of the relevant results. Subgrid techniques are possibly the greatest source of uncertainty in contemporary hydrodynamic simulations, so
a great deal of study in this direction is needed in order to produce a priori models that can satisfactorily explain the physics of galaxy formation.

1.3 Star formation in galaxies

1.3.1 The stellar content of galaxies

As far as observations are concerned, perhaps the most defining feature of a galaxy is its stellar content. Consequently, it is very important that a given theory of galaxy formation is able to adequately explain how stars form, how many stars form in a given environment and at what rate, as well as properties of the resulting distribution of stars, such as kinematics and chemistry. As we have discussed in the previous sections, dark matter dominates structure formation; baryonic gas gathers in dark matter haloes. However, the gas may be able to radiatively cool on a timescale that is shorter than the age of the halo. As it cools, it will lose pressure support and flow into the centre of the halo, eventually becoming self-gravitating as its density increases. This may lead to gas gathering in cold, dense clouds in which stars can be formed. A reasonable hypothesis for estimating the stellar content of a dark matter halo is to assume that it is some fraction of the available baryonic mass in the halo, which in turn can be expressed in terms of the cosmic baryon to dark matter ratio, or baryon fraction, $f_b = \Omega_b/\Omega_m$.

This is the basis behind a common method used to investigate the link between dark matter haloes and luminous galaxies, known as abundance matching (see e.g. Kravtsov et al. 2004a; Conroy et al. 2006; Behroozi et al. 2010; Behroozi et al. 2013; Moster et al. 2010, 2018). The basic assumption is that there is a one-to-one correlation between halo mass and stellar mass, such that the most massive galaxies live in the most massive haloes. It is therefore possible to take the halo mass function, as measured in cosmological N-body simulations and compare it to the observed stellar mass (or luminosity) function, allowing a prediction to be made about the stellar content of a given dark matter halo. There are several complications which can introduce uncertainties into this technique. The estimation of the halo mass function can be hampered at the low mass end by lack of resolution in the simulations, while the upper end suffers from small number statistics due to cosmic variance unless the volume simulated is large enough. Additionally, scatter arises due to mergers, tidal stripping and general stochasticity. The stellar mass side of the analysis is also obviously affected by observational uncertainties. However, the technique has been refined over some time and produces instructive results that are
Fig. 1.4 Integrated baryon conversion efficiency at $z = 0$ from abundance matching relations (Behroozi et al. 2013; Moster et al. 2018). Haloes at $\sim 10^{12} M_\odot$ are most efficient at forming stars. Efficiency drops with increasing or decreasing mass. The two abundance matching relations shown agree well at the peak of the relation, but begin to disagree at lower halo masses where the relation is not as well constrained. Also consistent with relations from complementary techniques such as halo occupation distributions (HOD) and conditional luminosity functions (see e.g. Guo et al. 2016, and references therein). Fig 1.4 shows results from two abundance matching relations, expressed in terms of an integrated baryon conversion efficiency. This is the stellar mass to (total) halo mass ratio, as a function of halo mass, normalised by the cosmic baryon fraction. It therefore indicates what fraction of the theoretically available baryon content of the halo has been converted into stars. Star formation efficiency (SFE) clearly varies with mass. Haloes around $10^{12} M_\odot$ are the most efficient, turning $\sim 20\%$ of the available baryons into stars, while efficiency drops considerably at higher and lower masses.

Observations of star formation in the Milky Way show that it occurs exclusively in giant molecular clouds (GMCs); given that the relevant scales in the ISM are much smaller than that of the host galaxy, it is a reasonable assumption to assume that the physics of star formation works in approximately the same manner in all galaxies. Thus, the overall SFE of a galaxy can be seen to be the product of the efficiency with which GMCs can form stars and the efficiency with which galaxies are able to form (and keep) GMCs. The latter is likely to be responsible
for the varying SFE with halo mass, but GMCs themselves are considerably less efficient than naive predictions would suggest. We will now examine both of these efficiencies in more detail.

1.3.2 The sites of star formation

In the Milky Way, \( \sim 20\% \) of the cold gas is in the form of H\(_2\) (the rest is dominated by neutral atomic hydrogen). It is observed to have an extremely clumpy distribution and the vast majority of H\(_2\) is found in GMCs, with masses of \( 10^5 - 10^6 \) M\(_\odot\) on scales of tens of parsecs. The main formation mechanism for H\(_2\) in the ISM occurs on the surface of dust grains where pairs of adsorbed hydrogen atoms recombine. The timescale for H\(_2\) formation on dust grains is

\[
t_{\text{form}} = 1.5 \times 10^7 \text{ yr} \left( \frac{n}{100 \text{ cm}^{-3}} \right)^{-1},
\]

assuming a typical Milky Way H\(_\text{I}\) to dust mass ratio of 100 (Hollenbach et al. 1971). In the absence of dust, H\(_2\) formation must occur through gas phase reactions which are many orders of magnitude slower. The main form of H\(_2\) destruction is photodissociation. In the unattenuated ISM radiation field, the lifetime of an H\(_2\) molecule is on the order of 600 yr (Stecher & Williams 1967). However, continuum attenuation by dust grains and line attenuation by the H\(_2\) itself (self-shielding) on the outskirts of the molecular cloud allow it to survive.

The classic model of the ISM is that of a three phase medium roughly in pressure equilibrium arising from thermal instability (see e.g. Field et al. 1969; McKee & Ostriker 1977). For a gas to be in thermal equilibrium, gas cooling and heating must balance. Because the net cooling rates are a function of the temperature and density of the gas (see section 1.2.3 for a brief overview), this leads to the possible thermal equilibrium states describing a complex path through \( \rho - T \) space. If we further assume that the medium has a constant pressure (which corresponds to \( \rho T = \text{const.} \)), then there are discrete locations in \( \rho - T \) space in which the gas can exist. Small isobaric\(^3\) perturbations of an initially static, homogeneous medium in thermal equilibrium will lead to gas accumulating at these possible equilibrium states. Any further perturbations will be restored to one of the equilibrium states. This leads to the prediction of an ISM consisting of a hot phase of \( T \sim 10^6 \) K, a warm phase of \( T \sim 10^4 \) K and a cold phase of \( T \sim 10^2 \) K, with a correspondingly

\(^3\)The assumption of the perturbations evolving in pressure equilibrium with the background is valid as long as the sound crossing time of perturbation is much shorter than the thermal timescale.
large variation in density across the three phases. This picture is consistent with the observed ISM.

However, observations of GMCs show that they are not in pressure equilibrium with the surrounding warm medium, with a pressure excess of roughly an order of magnitude (see the review by Burkert 2006, and references therein). This suggests that some other mechanisms are necessary to compress the gas into GMCs, as well as the likelihood that they are, to some extent, gravitationally bound objects. Gravitational instability, of the kind described in section 1.1.2 in the context of cosmic structure formation, is likely to contribute to the formation of GMCs. In a rotating galaxy the Jeans criterion is not the only requirement that must be met for gravitational instability. Conservation of angular momentum means that the Coriolis force acts as a restorative force against collapse in addition to pressure. A more appropriate metric is the Toomre parameter, $Q$, (Toomre 1964), where perturbations in a rotationally supported disc are unstable if

$$Q = \frac{c_s \kappa}{\pi G \Sigma} < 1,$$

where $\Sigma$ is the surface density and

$$\kappa = \sqrt{2} \left( \frac{V_{\text{circ}}^2}{R^2} + \frac{V_{\text{circ}}}{R} \frac{dV_{\text{circ}}}{dR} \right)^{1/2}$$

is the epicyclic frequency. For perturbations as small as GMCs to be unstable to gravitational collapse, $Q$ must be of the order of 0.1. In the warm phase of the ISM, the sound speed is too high ($c_s \sim 10 \text{ km s}^{-1}$) for such low values of $Q$. In the cold gas phase, meanwhile, ($c_s \sim 0.2 \text{ km s}^{-1}$) fragmentation is likely to occur at a wide range of scales. It is therefore possible that GMCs are formed by a process of a Toomre instability that is itself triggered by a thermal instability.

Several other mechanisms exist which may promote the growth of dense clouds on the GMC scale, often by triggering instabilities. The diffuse ISM is known to be highly turbulent across many scales (driven by gravitational, thermal and magnetodynamic instabilities, spiral density waves, SN explosions etc.). It is possible that GMCs are formed at the convergence points of large-scale turbulent flows (see e.g. Ballesteros-Paredes et al. 2007, and references therein), which may trigger thermal instability. Alternatively, while unable to form GMCs independently, the Parker instability (Parker 1966), which arises due the differential amplification of magnetic buoyancy in different areas of the diffuse ISM, may be able to induce gravitational instability in a near unstable disk. In addition to driving large-scale
turbulence, spiral density waves may promote GMC formation by inducing a local
density enhancement. Finally, galactic scale interactions, such as tidal forces from
or mergers with nearby galaxies can both induce turbulent motion in the ISM, as
well as causing gas to lose angular momentum and flow to the centre of the galaxy,
increasing density and triggering instabilities.

1.3.3 Star formation efficiency within giant molecular clouds

Having explored potential formation mechanisms of GMCs, we now turn our at-
tention to the formation of stars within them. A relationship between the gas
surface densities (of H\textsubscript{1} and H\textsubscript{2}) and star formation rate (SFR) surface densities
was proposed by Schmidt (1959) and extended by Kennicutt (1998). Data from
Kennicutt (1998) is shown in Fig. 1.5, showing globally averaged measurements for
normal spirals, the centres of those spirals and for starbursts. The data is well fit
by a single power law:

\[
\dot{\Sigma}_* = (2.5 \pm 0.7) \times 10^{-4} \left( \frac{\Sigma_{\text{gas}}}{M_\odot \text{pc}^{-2}} \right)^{1.4 \pm 0.15} M_\odot \text{yr}^{-1} \text{kpc}^{-2}.
\]  

(1.41)

Observations at lower surface densities from Bigiel et al. (2008) and Wyder et al.
(2009), as well some of the Kennicutt (1998) measurements, also shown in Fig. 1.5,
suggest a steepening of the relation below \( \sim 10 \, M_\odot \text{pc}^{-2} \). This is often interpreted
to indicate the presence of a threshold surface density for star formation, although
the potential causes are heavily debated. It has been suggested that the low surface
densities limit star formation by suppressing gravitational and/or shear instabilities
in the disc or by preventing the formation of H\textsubscript{2}. For the purposes of this summary,
we mainly consider the gas that is fit by a single power-law.

Under the assumption that the scale height of the disc is roughly constant
throughout the measurement, eq. 1.41 is equivalent to volume densities as well as
surface densities (with different normalisations of course). The slope of 1.4 is close
enough to suggest that the SFR is controlled by the self-gravity of the gas:

\[
\dot{\rho}_* = \epsilon_{\text{SF}} \frac{\rho_{\text{gas}}}{\tau_{\text{ff}}} \propto \rho_{\text{gas}}^{1.5}.
\]  

(1.42)

However, the normalisation factor, \( \epsilon_{\text{SF}} \) (which can be interpreted as the ratio of
the gas free-fall time to the gas consumption time), must be significantly less than unity
in order to provide the observed normalisation (on the order of 1%, see Krumholz
& Tan 2007, and references therein). This suggests that either a small fraction of
the GMC mass actual participates in the star formation or that the star formation
time scale is actually not the free-fall time, but rather $\tau_{ff}/\epsilon_{SF}$, which would imply additional physics.

It is also worth mentioning that in addition to the Schmidt law, $\dot{\Sigma}_* \propto \Sigma_{gas}^N$, the data are also well fit by $\dot{\Sigma}_* \propto \Sigma_{gas}/\tau_{dyn} \propto \Sigma_{gas}\Omega_{gas}$, where $\tau_{dyn}$ is the dynamic timescale of the disc at the outer edge of the star forming region or equivalently $\Omega_{gas}$ is the angular rotational velocity. Using this fit, the implication is that roughly 10% of the gas reservoir is consumed per orbital time. This can be explained by models in which the star formation is self-regulated to maintain a Toomre-stable disc (Silk 1997), is determined by the rate of collisions between molecular clouds (Tan 2000) or is triggered by spiral structure (Wyse & Silk 1989). However, while the trend of increasing SFE with increasing $\Omega_{gas}$ exists between galaxies, spatially resolved measurements reveal that it does not hold within galaxies (Leroy et al. 2008).

Likewise, spatially resolved measurements reveal that a universal single power
law, while working for global galaxy measurements, is not a good fit to the local relation (i.e. between annuli at different radii through a disc), with best fit values covering the range $N = 1 - 3$ (Bigiel et al. 2008). However, when only molecular gas surface density is used, a linear relation between gas and SFR surface density is found. The linear relation and normalisation found by Bigiel et al. (2008) imply a constant $\frac{\Sigma_*}{\Sigma_{\text{H}_2}} = (7 \pm 3) \times 10^{-10}$ yr$^{-1}$ corresponding to a gas consumption time of $\sim 1.5 \times 10^9$ yr. Assuming mean densities of GMCs of $100 - 500$ cm$^{-3}$, GMC free-fall times are $\sim 1 - 2$ Myr. This gives $\epsilon_{\text{SF,GMC}} = \tau_{\text{ff}} / \tau_{\text{SF}} \sim 0.1\%$. The ages of GMCs are known to be $\sim 10^7$ Myr (they are only correlated with stars younger than this); the low value of $\epsilon_{\text{SF,GMC}}$ and the fact that they are able to survive much longer than their free-fall times suggest that they are stabilised against global collapse.

GMCs have extremely complicated internal structures. As mentioned in the previous section, overall they typically extend over tens of parsecs with masses of $10^5 - 10^6$ M$_\odot$, corresponding to densities of $\sim 100 - 500$ cm$^{-3}$. Within their larger extent, however, a complex of filaments and clumps on all scales are observed. These range from clumps of $10^2 - 10^4$ M$_\odot$ on parsec scales ($n_{\text{H}_2} \sim 10^2 - 10^4$ cm$^{-3}$) to molecular cores of $0.1 - 10$ M$_\odot$ with $n_{\text{H}_2} > 10^5$ cm$^{-3}$. The formation of individual stars are thought to proceed from the collapse of molecular cores. Thus only the fraction of the GMC that is in the cores will become stars.

One possible source of pressure support that could prevent the global collapse of GMCs is magnetic fields. Considering the gravitational collapse of an isothermal non-rotating spherical cloud of mean density $\rho$ and radius $R$ embedded in a medium of constant pressure $P_0$ with a uniform magnetic field within the cloud of strength $B$ and ignoring magnetic surface terms, we can write the virial equation as (Spitzer 1968)

$$4\pi R^3 P_0 = 3\frac{M k_B T}{\mu} - \frac{1}{R} \left( \frac{3}{5} G M^2 - \frac{1}{3} R^4 B^2 \right),$$

where $M = \frac{4}{3} \pi R^3 \rho$ is the mass of the cloud, $k_B$ is Boltzmann’s constant, $T$ is the temperature of the region and $\mu$ is the mean mass per particle. Expressing the cloud radius in terms of the cloud mass and density, the two terms inside the parentheses in eq. 1.43 give a critical mass above which gravitational collapse against magnetic repulsion is possible:

$$M_{\text{cr}} = \frac{5^{3/2} B^3}{48\pi^2 G^{3/2} \rho^2} \approx \left(4 \times 10^6 \text{ M}_\odot \right) \left( \frac{n}{1 \text{ cm}^{-3}} \right)^{-2} \left( \frac{B}{3 \mu \text{G}} \right)^{-3}.$$  

(1.44)

Clouds with masses below $M_{\text{cr}}$ are referred to as subcritical (i.e. magnetostatically stable) and supercritical if their mass is greater than $M_{\text{cr}}$ (i.e. magnetostatically
unstable). Given a typical Milky Way magnetic field strength of $3\mu G$, it can be seen that $M_{cr}$ is larger than most GMC masses assuming a typical average density of the Milky Way ISM on the order of $< 1\, \text{cm}^{-3}$. If the ionization is high enough, the field will be frozen to the matter, meaning that the magnetic flux $\Phi = \pi R^2 B$ must remain constant, so it can be seen that $M_{cr}$ also remains constant. The implication, therefore, is that subcritical clouds remain subcritical even when compressed i.e. for a cloud to fragment into GMCs, clumps, molecular cores and eventually form stars, it must lower its magnetic flux. The process of ambipolar diffusion can in theory provide a suitable mechanism to achieve this while also regulating star formation (see e.g. Mestel & Spitzer 1956; Spitzer 1968; Mouschovias 1976; Shu 1977, 1983). The neutral particles in the cloud are only coupled to the magnetic field via collisions with ionized particles. Once the ionization fraction is sufficiently low, the coupling of the neutral population to the field is weakened, allowing neutral particles to diffuse through the field, reducing the magnetic flux. The ambipolar diffusion timescale would therefore regulate the star formation. However, the timescale on the scale of GMCs is $\tau_{ad} \sim 10^8 \, \text{yr}$ (see e.g. Elmegreen 1979), which is inconsistent with the lifetime of GMCs ($\sim 10^7 \, \text{yr}$). Additionally, observations suggest that most GMCs are supercritical or at most only marginally subcritical (see e.g. Crutcher 1999; Bourke et al. 2001), meaning that magnetic fields are unable to regulate star formation in this manner (see also the review Mac Low & Klessen 2004). The action of ambipolar diffusion in reducing magnetic flux would also result in a higher mass-to-flux ratio between cores and surrounding material in clouds than is observed (see e.g. Crutcher et al. 2009; Mouschovias & Tassis 2009; Lunttila et al. 2009; Santos-Lima et al. 2010; Lazarian et al. 2012; Bertram et al. 2012).

Alternatively, the supersonic turbulence that is observed throughout the ISM could regulate star formation (see e.g. Zuckerman & Palmer 1974; Zuckerman & Evans 1974; Larson 1981; Solomon et al. 1987; Falgarone et al. 1992; Ossenkopf & Mac Low 2002; Heyer & Brunt 2004; Schneider et al. 2011; Roman-Duval et al. 2011). There are various proposed sources of ISM turbulence and it is likely that all contribute to some extent. For example expanding SNRs (Balsara et al. 2004; de Avillez & Breitschwerdt 2005; Tamburro et al. 2009), growing HII regions (McKee 1989; Krumholz et al. 2006; Gritschneder et al. 2009; Peters et al. 2010; Goldbaum et al. 2011), compression and shocks from galactic spiral shocks (Dobbs & Bonnell 2008; Dobbs et al. 2008; Elmegreen 2009), gravitational contraction (Hoyle 1953; Vázquez-Semadeni et al. 1998; Klessen & Hennebelle 2010; Elmegreen & Burkert 2010; Federrath et al. 2011) and feedback from young stellar objects such as jets.
Fig. 1.6 Left: The distribution of densities relative to the mean density for a turbulent medium with Mach numbers of 1, 10 and 100 in the absence of magnetic fields. The variation with turbulent forcing parameter is also shown with fully solenoidal ($b = 1/3$), mixed ($b = 0.4$) and fully compressive ($b = 1$) modes. When the forcing is more solenoidal, the log-normal distribution is narrower. Right: The variation of the star formation efficiency per free-fall time (labelled here as SFR$_{ff}$ but equivalent to $\epsilon_{SF}$ used elsewhere in the text) as a function of virial parameter and Mach number adapted from Federrath & Klessen (2012) using their multi-freefall Krumholz & McKee (2005) model. This assumes a turbulent forcing parameter of 0.4 and no magnetic fields. Star formation becomes more efficient with increasing Mach number and decreasing virial parameter.

and outflows (Norman & Silk 1980; Banerjee et al. 2007; Nakamura & Li 2008; Cunningham et al. 2009; Wang et al. 2010) may excite compressive modes, while galactic rotation and magnetorotational instabilities may provide solenoidal modes (see e.g. Piontek & Ostriker 2004, 2007; Tamburro et al. 2009). Solenoidal modes are also excited from compressive forcing by nonlinear interactions of colliding shock fronts, by baroclinitiy, rotation and shear, and by viscosity such that even turbulence driven by purely compressive forcing contains an even mix of solenoidal and compressive power (see e.g. Federrath & Klessen 2012, and references therein). Turbulence acts to stabilise clouds on large scales, but induces local compressions resulting in the observed structure of filaments and cores that are the progenitors of star formation. The PDF of gas density in an isothermal turbulent medium can be modelled as a log-normal distribution$^4$,

$$p(s) = \frac{1}{\sqrt{2\pi}\sigma_s^2}\exp\left(-\frac{(s-s_0)^2}{2\sigma_s^2}\right), \quad (1.45)$$

$^4$Once densities are high enough for substructures to start collapsing, the distribution gains a power-law tail at high density (see e.g. Klessen 2000; Kainulainen et al. 2009).
where \( s \equiv \ln(\rho/\rho_0) \), \( \sigma_s \) is the standard distribution of \( s \), \( \rho_0 \) is the mean density and \( s_0 \) is the mean logarithmic density. The reason for this shape is that the hierarchical distribution of density structures arises from a multiplicative process in which shocks are amplified via interaction with other shocks (see e.g. Vazquez-Semadeni 1994; Passot & Vázquez-Semadeni 1998). Being multiplicative in \( \rho \), it is additive in \( s \) resulting in a Gaussian distribution of \( s \). The mean logarithmic density is related to the standard deviation of the distribution as \( s_0 = -\sigma_s^2/2 \) (Vazquez-Semadeni 1994; Federrath et al. 2008).

In the absence of magnetic fields, the width of the distribution is dependent on the Mach number, \( \mathcal{M} \), and turbulent forcing parameter, \( b \) (which takes the values 1 for purely compressive turbulence, 1/3 for purely solenoidal turbulence and 0.4 for a mixture) as (see e.g. Molina et al. 2012)

\[
\sigma_s^2 = \ln \left( 1 + b^2 \mathcal{M}^2 \right).
\]

Example distributions of \( \rho \) are plotted in the left panel of Fig. 1.6. The distribution broadens with increasing Mach number and a higher proportion of solenoidal modes with respect to compressive modes. Magnetic fields will tend to damp density fluctuations (see e.g. Ostriker et al. 2001; Price et al. 2011). With the inclusion of magnetic fields that are correlated with the density as \( B \propto \rho^{1/2} \), Molina et al. (2012) derive the dependence of \( \sigma_s \) to be

\[
\sigma_s^2 = \ln \left( 1 + b^2 \mathcal{M}^2 \frac{\beta}{\beta + 1} \right),
\]

where \( \beta = P_{th}/P_{mag} \) is the ratio of thermal to magnetic pressure (as \( B \to 0 \), \( \beta \to \infty \), recovering eq. 1.46).

As mentioned earlier, turbulence is expected to provide pressure support on large scales and promote dense structures that are unstable to gravitational collapse. The SFR is then expected to be regulated by the scale at which this regime change occurs and the mass of structures below this scale (see e.g. Krumholz & McKee 2005; Padoan & Nordlund 2011; Hennebelle & Chabrier 2011). Defining the star formation efficiency per free-fall time at the mean density of the gas, \( \epsilon_{SF} \) such that

\[
\dot{M}_* = \epsilon_{SF} \frac{M_{\text{cloud}}}{t_{ff}(\rho_0)},
\]

(which is analogous to eq. 1.42), the efficiency can be seen to be equivalent the mass fraction of gas with density above some critical density \( s_{\text{crit}} \) at which collapse can set it. We can therefore obtain this efficiency by integrating the density PDF
(Federrath & Klessen 2012):

\[
\epsilon_{\text{SF}} = \frac{\epsilon_{\text{core}}}{\phi_t} \int_{s_{\text{crit}}}^{\infty} \frac{t_{\text{ff}}(\rho_0)}{t_{\text{ff}}(\rho)} \frac{\rho}{\rho_0} p(s) ds.
\] (1.49)

The weighting by \( t_{\text{ff}}(\rho_0)/t_{\text{ff}}(\rho) \) accounts for the fact that gas with different densities have different free-fall times, while \( \phi_t \) is of order unity and accounts for uncertainty in that ratio (for details see Krumholz & McKee 2005). \( \epsilon_{\text{core}} \approx 0.25 - 0.7 \) (see e.g. Matzner & McKee 2000; Alves et al. 2007; André et al. 2010; Ward et al. 2012) accounts for the fact that not all gas that makes into molecular cores will accrete onto stars because of protostellar jets/winds/outflows. The value of \( s_{\text{crit}} \) adopted is dependent on the proposed criteria that determines the scales on which collapse sets in. For example, Krumholz & McKee (2005) defines this in terms of the Jeans length and the sonic scale (the scale at which the turbulence becomes subsonic, below which the global turbulent supersonic support is insignificant (see e.g. Vázquez-Semadeni et al. 2003; Mac Low & Klessen 2004; Federrath et al. 2010)) while Padoan & Nordlund (2011) defines it in terms of the magnetic shock jump conditions and the magnetic critical mass. Regardless, these models can all be expressed such that the dependence of \( \epsilon_{\text{SF}} \) is on the cloud-scale values of \( M, b, \beta \) and the virial parameter, \( \alpha_{\text{vir}} \), which is the ratio of twice the kinetic energy to the gravitational energy. The right panel of Fig. 1.6 (adapted from Federrath & Klessen 2012) shows the dependence of \( \epsilon_{\text{SF}} \) (SFR in their notation) on the virial parameter and the turbulent Mach number in the absence of magnetic fields (and with \( \epsilon_{\text{core}}/\phi_t = 1 \)). The efficiency drops sharply within increasing \( \alpha_{\text{vir}} \) and decreasing \( M \). Federrath & Klessen (2012) find that these types of models are good fits to numerical simulations (using sink particles for star formation). Given that the models provide a link between cloud scale variables and star formation efficiencies, they can potentially be used as the basis of subgrid schemes for star formation in galaxy-scale simulations (as in Kimm et al. 2017).

1.3.4 Feedback processes as regulators of star formation

Another major contributor to the regularisation of star formation comes from the stars themselves. As mentioned briefly in the previous section, protostellar winds are likely to reduce the efficiency of accretion onto the protostellar cores. However, the most dominant forms of stellar feedback on the whole GMC are produced by massive (OB) stars in the form of radiation (photoionization, photoelectric heating, radiation pressure), stellar winds and SNe. We leave a more detailed description of these processes until the introduction of Chapter 2, but we briefly mention a
few key points here. It can be seen that the typical lifetime of GMCs, $\sim 10^7$ yr, is comparable to that of massive stars, suggesting that the stars themselves are responsible for the destruction of their birth clouds. Expanding HII regions created around OB stars will gradually photoevaporate clouds. Stellar winds are likely to be subdominant in their impact on the cloud, but can alter the morphology of HII regions by clearing dense gas from around the stars at early stages (see e.g. Dale et al. 2014, Fig. 1.7 is adapted from this work and demonstrates the impact of photoionization and stellar winds on a GMC). Eventually, after $\sim 3-40$ Myr, the massive stars will explode as Type II SNe. These are likely to disperse their birth clouds, but the porous nature of the GMC can potentially lead to outflows travelling preferentially down low density channels while leaving dense clumps intact. In addition, the momentum coupled to the local gas by a SN is dependent on the local density, so the actions of stellar feedback are likely to influence the subsequent evolution of the SNR. More details and a review of the literature on these effects can be found in subsequent chapters.

While stellar feedback is able to influence the evolution (and lifespan) of the clouds that formed the stars, it can also have a major impact on larger scale galaxy

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5While Type Ia SNe can impact the chemistry of galaxies (most notably through their contribution to the relative abundances of alpha elements to iron), their rates are considerably lower than Type II SNe so do not contribute substantially to the feedback budget.
properties; this is a major theme of this thesis. Again, a more detailed analysis of the role feedback plays in regulating galaxy formation can be found in subsequent chapters, but we list the key points here. SNe in particular can have a major impact on galaxy evolution by driving galactic outflows which remove potentially star forming gas from the galaxy and may suppress the cosmic accretion. By impacting the distribution of gas in the galaxy, this will in turn influence a wide range of galaxy properties including morphologies, dark matter density profiles, kinematics, metal distribution (in particular enrichment of the CGM), escape fraction of ionizing photons (particularly relevant to cosmic reionization) to name a few. Aside from outflows, stellar feedback may act to heat the ISM, increasing the velocity dispersion and stabilising it against Toomre instabilities. As mentioned previously, feedback may contribute to the generation of ISM turbulence. SNRs are also thought to be a source of cosmic rays (through second order Fermi acceleration across the shock front) which can potentially drive outflows as well as providing a source of heating.

The relative ability of feedback to regulate star formation is also heavily dependent on the properties of system as a whole, most notably halo mass. For example, larger halo masses result in deeper potential wells that outflows have to climb. It appears that the decreasing efficiency of SNe to regulate star formation with increasing halo mass is responsible for the dependence of the integrated baryon conversion efficiency on halo mass at the faint end, as in Fig. 1.4 (or equivalently, the shape of the stellar mass function). At higher halo masses, regulation is believed to be provided in the form of feedback from AGN.
2. Supernova feedback in isolated simulations of galaxy formation

Separating physics from numerics

While feedback from massive stars exploding as supernovae (SNe) is thought to be one of the key ingredients regulating galaxy formation, theoretically it is still unclear how the available energy couples to the interstellar medium and how galactic scale outflows are launched. We present a novel implementation of six sub-grid SN feedback schemes in the moving-mesh code AREPO, including injections of thermal and/or kinetic energy, two parametrizations of delayed cooling feedback and a ‘mechanical’ feedback scheme that injects the correct amount of momentum depending on the relevant scale of the SN remnant resolved. All schemes make use of individually time-resolved SN events. Adopting isolated disc galaxy setups at different resolutions, with the highest resolution runs reasonably resolving the Sedov-Taylor phase of the SN, we aim to find a physically motivated scheme with as few tunable parameters as possible. As expected, simple injections of energy overcool at all but the highest resolution. Our delayed cooling schemes result in overstrong feedback, destroying the disc. The mechanical feedback scheme is efficient at suppressing star formation, agrees well with the Kennicutt-Schmidt relation and leads to converged star formation rates and galaxy morphologies with increasing resolution without fine tuning any parameters. However, we find it difficult to produce outflows with high enough mass loading factors at all but the highest resolution, indicating either that we have oversimplified the evolution of unresolved SN remnants, require other stellar feedback processes to be included, require a better star formation prescription or most likely some combination of these issues.

2.1 Introduction

In the \( \Lambda \)CDM model of cosmology, dark matter dominates large scale structure formation. Gas gathers in the potential wells of dark matter haloes, where it may radiatively cool and hence form stars. This baryonic matter makes up the visible component of galaxies. This picture alone is not sufficient to reproduce observations. A naive determination of the expected star formation rate (SFR) based on
a typical dynamical time yields excessive values. In fact, star formation occurs on much longer timescales of the order of 20 - 100 dynamical times and has an efficiency of only a few percent (see for example Zuckerman & Palmer 1974; Williams & McKee 1997; Kennicutt 1998; Evans 1999; Krumholz & Tan 2007; Evans et al. 2009). Thus, some form of feedback process or processes need to be invoked to explain this discrepancy. At high halo masses, this may be provided by an active galactic nucleus (AGN), but at lower masses stellar feedback dominates, mainly from high mass stars in the form of stellar winds, supernovae (SNe), photoionization and radiation pressure.

It is worth emphasising that it is not enough merely to halt the conversion of gas to stars as some fraction of the accreted mass must be ejected out of the galaxy. Without strong feedback, the baryon fraction of galaxy models are far in excess of observations (e.g. White & Frenk 1991; Kereš et al. 2009). In addition, the observed circumgalactic medium (CGM) is enriched with metals, requiring baryons to have made it out from sites of star formation embedded within the galaxies themselves (e.g. Aguirre et al. 2001; Pettini et al. 2003; Songaila 2005, 2006; Martin et al. 2010). Such outflows are observed, moving at hundreds of km s\(^{-1}\) (see for example the review by Veilleux et al. 2005). Observations suggest that the ratio of mass outflow rate to SFR (i.e. the mass loading factor) must be at least unity or above (see e.g. Bland-Hawthorn et al. 2007; Schroetter et al. 2015). This is borne out by theoretical models (Oppenheimer & Davé 2006; Sales et al. 2010; Genel et al. 2012; Shen et al. 2012; Davé et al. 2013; Puchwein & Springel 2013; Vogelsberger et al. 2013; Hopkins et al. 2014; Mitra et al. 2015; Christensen et al. 2016).

While the observational evidence for SFR regularisation and outflow driving is manifest, precisely how these mechanisms operate is as yet unclear. Here numerical hydrodynamic simulations of galaxy formation are useful tools. Unfortunately, the scales on which stellar feedback operates (parsecs and below) are many orders of magnitude below the characteristic scales of galaxies and the surrounding CGM we wish to simulate.

This ideally needed dynamic range is beyond the reach of current state-of-the-art simulations, requiring the representation of the effects of unresolved processes by adopting so-called ‘sub-grid’ schemes. For large scale cosmological simulations, where the interstellar medium (ISM) is poorly resolved, these schemes must rely on dealing with stellar feedback at a high level of abstraction. For example, such approaches may use effective equations of state to approximate the effect of a multiphase ISM pressurised by feedback energy (e.g. Springel & Hernquist 2003; Teyssier et al. 2010). Winds are often added with some predetermined mass loading,
either temporarily decoupling outflowing gas from the hydrodynamics, imposing some minimum threshold temperature of the wind ejecta or switching off radiative cooling losses for a given amount of time, to ensure sufficiently strong driving (e.g. Springel & Hernquist 2003; Oppenheimer & Davé 2006; Dalla Vecchia & Schaye 2008; Sales et al. 2010). Such schemes are presently necessary to model large samples of galaxies but lack predictive power on small scales. However, if the target of a simulation is a single galaxy, either in an idealised, isolated setup or in a cosmological ‘zoom-in’, then the higher resolution available enables the adoption of more explicit models of feedback, allowing investigations of how feedback arises on comparatively smaller scales to be carried out.

Nevertheless, even in individual galaxy simulations the resolution requirements are still severe. In the case of SNe, one of the main obstacles to physically consistent coupling of SN energy to the ISM is the ability to resolve the Sedov-Taylor phase of a SN remnant. The expansion of SN remnants has been well studied and can be broken down into several distinct regimes (Woltjer 1972). The SN explosion ejects material into the ISM with typical kinetic energies of $10^{51}$ ergs. The SN ejecta expands relatively unhindered into the ISM as long as the mass swept up in the forward shock is smaller than the ejecta mass. Concurrently, the reverse shock heats up the gas inside the remnants, leading to high temperatures and pressures. Radiative losses are negligible so the expansion proceeds adiabatically into the surrounding medium, which marks the Sedov-Taylor phase. During this phase, the momentum of the remnant is boosted by up to an order of magnitude (Taylor 1950; Sedov 1959; Chevalier 1974; Cioffi et al. 1988; Blondin et al. 1998; Kim & Ostriker 2015; Martizzi et al. 2015). Eventually, a thin, dense shell builds up at the shock front and radiative losses become important, triggering the transition from energy conserving to momentum conserving evolution. Because of the large increase in momentum that occurs during the adiabatic expansion, merely injecting energy (whether thermal or kinetic) into the surrounding gas without properly resolving the length scales corresponding to the Sedov-Taylor phase results in a severe underestimation of the amount of momentum imparted to the ISM. Kim & Ostriker (2015) found that the minimum requirements for correctly modelling the evolution of SNe in this manner are that the shell forming radius, $r_{SF}$, is resolved by three resolution elements. For evolution in an inhomogeneous medium, they quantified that $r_{SF} = 30 \text{ pc} \left(\frac{n}{\text{cm}^{-3}}\right)^{-0.46}$, meaning that at a density of $100 \text{ cm}^{-3}$ (typical mean density for a giant molecular cloud) the resolution requirement is $\sim 1 \text{ pc}$. Failure to meet these requirements when using a simple injection of SN energy will result in ‘overcooling’ as the energy is radiated away before it can do
Many strategies to circumvent this issue exist in the literature. One implicit solution is to inject the energy of several SNe simultaneously resulting in more energetic explosions. Often, this is achieved simply by injecting a star particle’s entire feedback energy budget at once, either instantaneously or after some pre-determined delay time. The strength of this effect is therefore tied to the star particle mass. Alternatively, a stochastic feedback approach, such as that proposed in Dalla Vecchia & Schaye (2012), may be adopted, in which SN energy is redistributed in time and space to produce fewer, more energetic events guaranteeing the overcooling problem is avoided. Such schemes conserve the total feedback energy in a globally averaged sense, but lose the connection to individual SN events and are not spatially consistent. If the simulation is of a coarse resolution and the structure of the ISM is not resolvable/of interest, this may be an acceptable compromise.

A different class of approaches involves switching off the radiative cooling of gas that has received feedback energy, enforcing an adiabatic phase, for some length of time (see e.g. Stinson et al. 2006; Governato et al. 2010; Agertz et al. 2011; Teyssier et al. 2013). The length of time by which cooling is delayed is somewhat of a tunable parameter, particularly in simulations with coarse resolution, but physically motivated parameters can be arrived at by analytical arguments (see for example the appendix of Dubois et al. 2015). A downside of ‘delayed cooling’ models is that the radiative cooling of the gas is physically correct, even if the resolution effects responsible for the overcooling phenomena are not. Thus it is possible for gas to occupy unphysical regions of temperature-density phase diagrams when it should have cooled.

In an alternative to the ‘delayed cooling’ schemes, it is possible to take account of the momentum boost in the missed adiabatic phase rather than enforcing such a phase. Some schemes skip the Sedov-Taylor phase entirely, putting in a bubble at some fixed radius and adjusting the kinetic energy of the gas inside to match the analytically determined values assuming some mass loading (see e.g. Dubois & Teyssier 2008). Others determine the stage of a remnant’s evolution that can be resolved and boost the momentum by some appropriate factor determined either analytically (e.g. Hopkins et al. 2014; Kimm & Cen 2014) or by making use of fits to high resolution simulations of SN remnant evolution (e.g. Martizzi et al. 2015 as employed in Martizzi et al. 2016). These schemes are often referred to as mechanical feedback. They feature few (if any) explicitly tunable parameters, but rely on assumptions about the structure of the ISM at small scales and how the remnant will interact with it (e.g. models typically assume a spherically symmet-
Supernova feedback in isolated simulations of galaxy formation

...ric evolution in an homogeneous background medium). A porous ISM structure caused by turbulence may allow the remnant to propagate preferentially down low density channels (Iffrig & Hennebelle 2015; Kim & Ostriker 2015; Martizzi et al. 2015; Walch & Naab 2015; Li 2015; Haid et al. 2016) though the net effect of this phenomenon is not well constrained and possibly introduces further free parameters into the model.

Of course, SNe are not the only form of stellar feedback. It is possible for photoionization to break up star forming clouds prior to the first SNe occurring (Vázquez-Semadeni et al. 2010; Walch et al. 2012; Dale et al. 2014; Sales et al. 2014). Winds from massive stars are unable to completely disrupt $10^4 - 10^5 M_\odot$ clouds, but can carve cavities of $\sim 10$ pc which may enhance subsequent SNe feedback (Dale et al. 2014). Radiation pressure can in principle supply as much momentum as stellar winds (see e.g. Leitherer et al. 1999), though it is difficult to assess the extent to which this can be coupled to the ISM. On the one hand, HII regions created by massive stars will blunt the impact of radiation pressure, rendering the ISM transparent to Lyman-limit photons, but in the presence of dust, multiple scattering of IR photons can boost the momentum input to the ISM by up to a few orders of magnitude (Murray et al. 2010). Using sub-grid models of radiation pressure feedback it has been found that boost factors of $\sim 10 - 100$ are necessary to drive strong outflows (Hopkins et al. 2011, 2012a,b; Agertz et al. 2013; Aumer et al. 2013; Roškar et al. 2014; Agertz & Kravtsov 2015). However, using full radiative hydrodynamics (RHD), Rosdahl et al. (2015) concluded that radiation pressure is unable to drive strong outflows in their simulations, although they are unable to resolve gas at high enough densities to become significantly optically thick to IR photons. Nevertheless, a simple boosting of the IR optical depths resulted in suppressing star formation and smoothing of the disc without generating outflows. In reality, all of these stellar feedback mechanisms will interact in a complex manner. For example, the FIRE project (Hopkins et al. 2014) has produced encouraging results by including multiple stellar feedback processes in sub-grid fashion, creating realistic looking galaxies relative to observations. However, it is clear that before trying to unpick the interaction of different processes and their impact on galaxy formation, it is crucial to understand the numerical consequences of the individual feedback schemes.

To this end, in this work, we carry out a detailed study of various flavours of SN feedback prescriptions commonly found in the literature. We perform simulations of idealised, isolated galaxy models, in the absence of other feedback prescriptions and with a simple star formation law, in order to provide as clean a comparison...
as possible. The schemes tested are all chosen to work with individually time resolved SN events, providing as direct a link to the locations and timescales of star formation as possible (e.g. we do not consider stochastic feedback such as Dalla Vecchia & Schaye 2012), and are optimised for isolated or cosmological zoom-in simulations (rather than cosmological boxes). We carry out our fiducial simulations of a $10^{10}$ M$_\odot$ system at three resolutions (the highest of which is chosen to largely eliminate the overcooling problem in a simple thermal dump scheme) in order to test convergence properties, trialing six sub-grid feedback schemes. Having presented our main findings with respect to resulting galaxy morphologies, SFRs, and outflow properties as function of feedback scheme, we briefly examine how these results depend on the mass of the galaxy and simple changes of the star formation prescription.

### 2.2 Methodology

#### 2.2.1 Basic code setup

We make use of the moving-mesh code AREPO (Springel 2010) with our own novel implementation of star formation and SN feedback (described below). AREPO uses a quasi-Lagrangian finite volume technique, solving hydrodynamics on an unstructured mesh determined by a Voronoi tessellation of discrete mesh-generating points. These points move with the local gas velocity (with the addition of minor corrections to allow for cell regularisation). By moving the mesh with the fluid and employing a smoothly varying refinement and derefinement scheme, AREPO is able to keep cell masses constant (to within a factor $\sim 2$). AREPO benefits from many of the advantages inherent to traditional Lagrangian approaches (e.g. smoothed particle hydrodynamics (SPH)), such as continuously varying resolution with density and Galilean invariance, while retaining advantages of contemporary Eulerian codes (i.e. adaptive mesh refinement (AMR)) such as more accurate resolution of shocks, contact discontinuities and fluid instabilities (Bauer & Springel 2012; Kereš et al. 2012; Sijacki et al. 2012; Torrey et al. 2012; Vogelsberger et al. 2012). We include radiative cooling from both primordial species and metal-lines as presented in Vogelsberger et al. (2013): primordial heating and cooling rates are calculated using cooling, recombination and collisional rates from Cen (1992) and Katz et al. (1996), while lookup tables pre-calculated with the photoionization code CLOUDY are used to obtain the metal cooling rates. Note that in this work we do not include a UV background.
2.2.2 Non-thermal pressure floor

Failing to sufficiently resolve the Jeans length can result in artificial fragmentation (Truelove et al. 1997). To avoid this we include a non-thermal pressure floor to ensure that the Jeans length is resolved by \( N_J \) cells, i.e.

\[
P_{\text{min}} = \frac{N_J^2 \Delta x^2 G \rho^2}{\pi \gamma},
\]

where \( \Delta x \) is the cell diameter, \( \rho \) is the gas density and \( \gamma = 5/3 \) is the adiabatic index. In principle, at sufficiently high resolution if feedback is able to entirely prevent gas from entering a phase where it is vulnerable to artificial fragmentation, it may be possible to avoid the use of a pressure floor. This would ideally prevent the risk of suppressing physical fragmentation which may occur when a pressure floor is in place. Alternatively, the star formation prescription adopted could be formulated to ensure that gas is turned into stars before artificial fragmentation occurs. However, as we only include SN feedback (note that there is a delay of \( \sim 3 \) Myr before the first SNe go off) and wish to study the effects of the feedback without a more involved method of modelling star formation (see below), we use a pressure floor to ensure numerically meaningful gas conditions prior to feedback.

Various values for \( N_J \) can be found in the literature. We find that the value required varies depending on choice of code, cooling prescriptions, initial conditions, resolution and included sub-grid physics. The choice is therefore somewhat arbitrary and often ill defined. By performing an array of numerical experiments, we find that \( N_J = 8 \) is a reasonable choice for 1000 M\(_\odot\) cell resolution (see Fig. 2.16 in Appendix 2.A). It should be noted that in the absence of feedback, the choice of \( N_J \) has a significant impact on the total stellar mass formed (see Fig. 2.17). For the purposes of the comparison of feedback implementations in this work, our choice is therefore motivated by our requirements to avoid the opposite extremes of artificial fragmentation or total suppression of star formation by the pressure floor.

Using a fixed value of \( N_J \) with different resolutions ensures that the Jeans length is always resolved by the same number of cells. This means that, by design, fragmentation is allowed to occur on smaller scales as simulations move to higher resolutions and the minimum resolvable scale decreases. While under most circumstances this is a desirable behaviour, the resulting lack of convergence in the absence of feedback makes a meaningful study of the resolution dependence of SN feedback schemes impossible. Thus, for this work, we adopt the scaling \( N_J = 8(m_{\text{cell}}/1000 \ M_\odot)^{-1/3} \) such that the pressure floor corresponds to resolving
the same length-scale across all resolutions\(^1\). This results in relatively similar gas morphologies, temperature and density distributions and SFRs in the absence of feedback at all numerical resolutions explored (see Fig. 2.18 in Appendix 2.A). With this choice of the pressure floor scaling, starting with relatively similar disc properties in different resolution runs we can more readily isolate how feedback operates at different resolutions. A much more detailed discussion of the use of the pressure floor in this work and its effects on the simulations is presented in Appendix 2.A.

### 2.2.3 Star formation

In our model, gas is marked as star forming if it is above some density threshold \(n_{\text{SF}}\). We then compute a star formation rate for the gas based on a simple Schmidt law, using the almost ubiquitous expression

\[
\dot{\rho}_* = \epsilon_{\text{SF}} \frac{\rho}{t_{\text{ff}}},
\]

(2.2)

where \(\rho\) is the gas density, \(\epsilon_{\text{SF}}\) is some efficiency and \(t_{\text{ff}} = \sqrt{3\pi/32G\rho}\) is the free-fall time. We use a fiducial value of \(n_{\text{SF}} = 10\ \text{cm}^{-3}\) and \(\epsilon_{\text{SF}} = 1.5\%\) (chosen to match observed efficiencies in dense gas, see e.g. Krumholz & Tan 2007, and references therein). These values are kept the same across all resolutions for our fiducial simulations (they are an appropriate choice for all resolutions explored), with the aim of removing the dependence the choice of star formation law prescription and allowing us to assess the convergence properties of the SN schemes alone\(^2\). We then use these rates to stochastically convert gas cells to star particles (representing a single stellar population).

### 2.2.4 Supernova feedback

Our implementation of SN feedback is directly related to individual star particles and discretely resolves individual SNe in time. This is in contrast to implementations which inject energy continuously at some rate related to the SFR and to methods in which a fixed quantity of energy per stellar mass in injected, possibly after some delay. Injecting the energy of multiple SNe at once will help avoid the overcooling problem (the radius of the remnant at the end of the Sedov-Taylor

\(^1\)For gas at 1000 K, \(N_J = 8\) means that structures smaller than \(\sim 50\ \text{pc}\) are suppressed.

\(^2\)However, see Section 2.3.8 where we present results with a higher density threshold value of \(n_{\text{SF}} = 100\ \text{cm}^{-3}\) or a higher efficiency of \(\epsilon_{\text{SF}} = 15\%\).
phase has a dependence on injected energy as $E^{0.29}$ (Kim & Ostriker 2015)). However, the local evolution of the ISM with time as it evolves prior to the first SNe and as SNe occur sequentially (for example enhancing the strength of subsequent SNe) is non-trivial. Failing to resolve individual SNe in time potentially misses important physics. Therefore, each timestep, for each star particle, we tabulate SN rates, $\dot{N}_{SN}$, as a function of age and metallicity from Starburst99 (Leitherer et al. 1999) assuming a Kroupa (2002) IMF. We then draw the number of SNe that occur from a Poisson distribution with a mean $\bar{N}_{SN} = \dot{N}_{SN}\Delta t$, where $\Delta t$ is the timestep. We further impose a timestep limiter for star particles such that $\bar{N}_{SN} \ll 1$ to ensure that SNe are individually resolved in time.

When a SN occurs, mass, metals, energy and/or momentum (depending on the feedback scheme, see below) are deposited into the gas cell hosting the star particle and its immediate neighbours (i.e. all cells that share a face with the host cell). Fig. 2.1 demonstrates the procedure schematically. The various quantities are distributed amongst these cells using a weighting scheme that aims to guarantee an isotropic distribution. This contrasts with the SPH-like (mass, volume etc.) weighting schemes commonly used in Lagrangian codes. Because higher density regions will contain more resolution elements, such a weighting scheme will preferentially inject feedback quantities perpendicular to the local density gradient. In the worst case scenario, we have found that this manifests itself in the unphysical driving of strong feedback ‘rings’ through the plane of a thin disc, similar to those reported in Hopkins et al. (2017) and Hopkins et al. (2018) (see our Appendix 2.B for more details). Our weighting scheme is based on the ‘vector weighting’ scheme from Hopkins et al. (2017). Essentially, the quantities are weighted both by the solid angle subtended by the adjoining cell face and by projection operators to enforce isotropy. To compute these quantities we use the mesh geometry used in the hydrodynamic calculation. For reasons of numerical simplification we take the centre of the SNe to be the mesh generating point of the cell hosting the star particle, rather than the star particle itself. The effects of this are small, since by definition the star particle is spatially unresolved in the context of the hydrodynamic resolution. Further simplifications to the Hopkins et al. (2017) scheme arise since the Voronoi tessellation guarantees that cell face norms are aligned with the position vector between two mesh generating points and that cell faces lie exactly halfway between the two mesh generating points in that direction. Note that the centre of the cell face is not guaranteed to lie on the line between two mesh generating points. However, we take this as an approximation. Due to the mesh regularisation schemes used in Arepo, we find this to be reasonable.
Fig. 2.1 A schematic illustrating the distribution of mass, metals, energy and momentum from the star star particle to the local gas. First, the cell hosting the star particle (red star) is found by searching for the nearest mesh generating point (green). The host cell receives a fraction, $f_{\text{host}}$, of the feedback quantities while the rest is distributed to the neighbouring cells. For scalar quantities, the distribution is effectively weighted by the solid angle subtended by the intervening cell face, as resolved from the host cell mesh generating point, under the assumption that this can be approximated as a cone configuration (eq. 2.6), indicated by the dotted lines. However, merely weighting by the solid angle does not guarantee that the vector sum of the momenta is zero, so an additional correction (eq. 2.5) is required. The black arrows on the schematic indicate the vectors between the mesh generating points of the host cell and the neighbouring cells (blue); these are the $x_j$ used in eq. 2.4–2.8.

We first find the cell that contains the star particle, hereafter referred to as the host cell. For each of the neighbour cells (cells that share a face with the host cell), $i$, we determine the vector weight $\bar{w}_i$ defined as

$$\bar{w}_i = \frac{w_i}{\sum_j |w_j|} (1 - f_{\text{host}}), \quad (2.3)$$

where the sum over $j$ is over all neighbour cells including $i$,

$$w_i = \omega_i \sum_{\alpha=\pm} \sum_\alpha \left(\hat{x}_i^\alpha\right)^\alpha f_{\pm}^\alpha, \quad (2.4)$$
\[
    f_{\pm}^\alpha = \left\{ \frac{1}{2} \left[ 1 + \left( \frac{1}{2} + \frac{1}{2} \right) \right] \right\}^{1/2},
    \tag{2.5}
\]

\[
    \omega_i = \frac{1}{2} \left( 1 - \frac{1}{\sqrt{1 + 4A_i/(\pi |x_i|^2)}} \right),
    \tag{2.6}
\]

\[
    (\hat{x}_i^+)^\alpha = |x_i|^{-1} \text{MAX} (x_i^\alpha, 0)_{\alpha=x,y,z},
    \tag{2.7}
\]

\[
    (\hat{x}_i^-)^\alpha = |x_i|^{-1} \text{MIN} (x_i^\alpha, 0)_{\alpha=x,y,z},
    \tag{2.8}
\]

where \( f_{\text{host}} \) is the fraction of feedback quantities given to the host cell, \( A_i \) is the area of the face between the neighbour and the host cell, \( x_i \) is the position vector between the mesh generating points of the neighbour to the host, the superscript \( \alpha \) denotes the component in a given Cartesian direction, \( x, y \) or \( z \), while the \( + \) and \( - \) denote components with either a positive or negative value respectively. Given a total ejecta mass \( m_{\text{ej}} \) and SN energy \( E_{SN} \), the total momentum to be injected in the rest frame of the star particle is

\[
    p_{\text{tot}} = \sqrt{2m_{\text{ej}}f_{\text{kin}}E_{SN}},
    \tag{2.9}
\]

where \( f_{\text{kin}} \) is the fraction of ejecta energy that is in kinetic form (which we vary throughout this work). The portion of mass, momentum and total energy each cell receives (in the rest frame of the star particle) is then

\[
    \Delta m_i = |\bar{w}_i| m_{\text{ej}} \tag{2.10}
\]

\[
    \Delta p_i = \bar{w}_i p_{\text{tot}} \tag{2.11}
\]

\[
    \Delta E_i = |\bar{w}_i| E_{SN}. \tag{2.12}
\]

Eq. 2.6 can be seen to be the fraction of the unit sphere subtended by the cell face, \( \omega_i = \Delta\Omega_i/4\pi \), approximating the configuration as a cone\(^3\). Weighting by solid angle alone is sufficient for the scalar quantities. For the momentum, a simple approach would be to weight the magnitude of the momentum by the solid angle and point the vector along \( x_i \). However, unless the neighbouring mesh generating points are exactly isotropically distributed, this does not guarantee \( \Sigma_j \Delta p_j = 0 \). We must therefore introduce the correction factor given by eq. 2.5 as used in eq. 2.4. We can derive eq. 2.5 by rewriting the vector state-
ment of momentum conservation in terms of the Cartesian components: using eq. 2.4 we have $\sum_j (\Delta p_j)\alpha = (p_{tot}/\sum_j \omega_j)[f^\alpha\sum_j \omega_j |\hat{x}_j|\alpha - f^\alpha\sum_j \omega_j |\hat{y}_j|\alpha] = 0$. Since $p_{tot}$ and $\sum_j \omega_j$ are positive-definite, the terms inside the square brackets must sum to zero. Thus, defining $g^\alpha_\pm = \sum_j \omega_j |\hat{x}_j|\alpha$, we have $f^\alpha_+ g^\alpha_+ = f^\alpha_- g^\alpha_-$. A second constraint comes from requiring that the sum of the magnitudes of the total momentum coupled must not be modified by the weighting. This can be expressed as $(f^\alpha_+ g^\alpha_+)^2 + (f^\alpha_- g^\alpha_-)^2 = (g^\alpha_+)^2 + (g^\alpha_-)^2$. Eq. 2.5 is the function that satisfies these constraints.

Transforming back to the simulation frame (i.e. the rest frame of the simulated volume), the momentum and energy fluxes become

$$\Delta p'_i = \Delta p_i + \Delta m_i v_* \quad (2.13)$$

$$\Delta E'_i = \Delta E_i + \frac{1}{2\Delta m_i} \left( |\Delta p'_i|^2 - |\Delta p_i|^2 \right) \quad (2.14)$$

where $v_*$ is the velocity of the star particle in the simulation frame. Note that this implicitly deals with any momentum cancellation, i.e. the ‘lost’ kinetic energy becomes thermal energy.

The host cell receives the following mass and energy

$$\Delta m_{host} = f_{host} m_{ej} \quad (2.15)$$

$$\Delta E_{host} = f_{host} E_{SN} + \frac{1}{2} f_{host} m_{ej} |v_* - v_{host}|^2 \quad (2.16)$$

The final term in equation (2.16) assumes complete thermalisation of the kinetic energy carried by the star particle. Empirically, we find that the mean number of neighbouring cells is $\sim 20$. We therefore adopt $f_{host} = 5\%$ to evenly distribute feedback quantities. In practice, we find a very weak dependence on the value of $f_{host}$. In simulations described as containing no feedback, the host cell and neighbours receive mass and metals as described above but their energy and momentum are not altered. We adopt $m_{ej} = 10 M_\odot$, of which 2 $M_\odot$ is in metals (i.e. a metallicity of 0.2), and $E_{SN} = 10^{51}$ ergs throughout this work.

**Classical feedback schemes**

For the purpose of this work, we refer to schemes that employ a simple dump of thermal and/or kinetic energy as classical feedback schemes. These use the methods outlined above with some value of $f_{kin}$. For pure thermal feedback, we use $f_{kin} = 0$. 


For pure kinetic feedback\textsuperscript{4}, we use \( f_{\text{kin}} = 1 \). We also trial a mixed feedback scheme that uses \( f_{\text{kin}} = 0.28 \), which distributes the energy into the ratio expected during the Sedov-Taylor phase (e.g. see Ostriker & McKee (1988); Cioffi et al. (1988) for analytical arguments, also see Kim & Ostriker (2015) for an example of this in a numerical simulation).

\textbf{Delayed Cooling}

Additionally to the classical feedback schemes we adopt a feedback prescription based on the delayed cooling method of Teyssier et al. (2013). This method aims to take into account (sub-grid) non-thermal processes that might store some of the feedback energy, such as, for example, unresolved turbulence, magnetic fields and cosmic rays. The timescales on which these processes dissipate energy is longer than the cooling time of the thermal component, so energy may be stored for longer and released gradually. We introduce a new variable \( u_{\text{FB}} \) which is used to record the energy density from feedback that gas particles currently possess and is advected with the gas flow, acting as a passive Lagrangian tracer (which is to say it is not directly involved in the hydrodynamics). When a gas cell is involved in a SN event, feedback energy is injected as described above with \( f_{\text{kin}} = 0 \) (i.e. entirely thermally apart from the momentum conserved from the star particle). The amount of energy received is also added to \( u_{\text{FB}} \). This feedback energy store is allowed to dissipate as

\[ \frac{du_{\text{FB}}}{dt} = -\frac{u_{\text{FB}}}{t_{\text{diss}}}, \]

where \( t_{\text{diss}} \) is some dissipation timescale as in Teyssier et al. (2013). Note that \( u_{\text{FB}} \) can also be increased if the gas cell is involved in another SN event. We compute an effective velocity dispersion corresponding to the feedback energy,

\[ \sigma_{\text{FB}} = \sqrt{2u_{\text{FB}}}, \]

and the gas particle is not allowed to cool if this velocity dispersion is above some threshold. Following Teyssier et al. (2013) we use

\[ \Lambda = 0 \text{ if } \sigma_{\text{FB}} > 10 \text{ km s}^{-1}. \]

\textsuperscript{4}This should not be confused with other schemes sometimes referred to as ‘kinetic’ that boost the momentum input by some fixed mass loading factor (see Dubois & Teyssier 2008; Kimm et al. 2015; Rosdahl et al. 2017).
The motivation for switching off cooling when $\sigma_{FB}$ is above this threshold is to mimic a non-thermal contribution to the pressure. Once the non-thermal contribution becomes comparable to the thermal contribution, cooling is allowed to continue as normal. We also trial a larger threshold value of 100 km s$^{-1}$ in Appendix 2.C.

We use a fixed value for the dissipation time of 10 Myr, as in Teyssier et al. (2013). We also trial a variable dissipation time, based on the effective crossing time for the turbulence within a cell,

$$t_{diss} = \frac{\Delta x}{\sigma_{FB}}, \quad (2.20)$$

where $\Delta x$ is the diameter of the cell.

**Mechanical feedback**

In this feedback scheme we aim to account for the $PdV$ work done during the Sedov-Taylor phase of the SNe remnant expansion, where the momentum can be boosted by around an order of magnitude. The correct momentum to couple to the ISM therefore depends on the stage of the expansion (alternatively parametrized in terms of swept up mass), limited by the final momentum at the point when the remnant exits the Sedov-Taylor phase. Several such schemes exist in the literature (see e.g. Hopkins et al. 2014, 2017; Hopkins et al. 2018; Kimm & Cen 2014; Kimm et al. 2015; Martizzi et al. 2015). In our mechanical feedback scheme, momentum calculated in equation (2.13) is enhanced as follows

$$\Delta p''_i = \Delta p'_i \text{MIN} \left[ \sqrt{1 + \frac{m_i}{\Delta m_i}}, \frac{p_{\text{fin}}}{p_{\text{tot}}} \right], \quad (2.21)$$

where $p_{\text{fin}}$ is the momentum as the remnant transitions to the snowplough phase (Blondin et al. 1998; Thornton et al. 1998; Geen et al. 2015; Kim & Ostriker 2015; Martizzi et al. 2015). Following Kimm et al. (2015) we adopt

$$p_{\text{fin}} = 3 \times 10^5 \text{ km s}^{-1} \text{ M}_\odot E_{51}^{16/17} n_{\text{SN}}^{-2/17} Z_{\text{SN}}^{-0.14}, \quad (2.22)$$

where $E_{51} = (E_{SN}/10^{51} \text{ ergs})$ is the energy of the SN (for our individually time-resolved SNe, $E_{51} \equiv 1$), while $n_{\text{SN}} = (n_H/\text{cm}^{-3})$ and $Z_{\text{SN}} = \text{MAX}(Z/Z_\odot, 0.01)$ are the hydrogen number density and metallicity of the ambient gas, respectively. Note that we calculate $\Delta p''_i$ for each cell involved in the SN event independently.
Table 2.1 Initial conditions of the three disc galaxies modelled in this work, referred to as ‘Small’, ‘Fiducial’ and ‘Large’. We list the total mass of the galaxy, $M_{\text{tot}}$, (excluding the CGM, which is negligible), the halo virial radius, $R_{\text{vir}}$, the mass in the disc component, $M_{\text{disc}}$, the fraction of the disc component in gas, $f_{\text{gas}}$, the scale radius of the disc, $r_s$, the scale height of the stellar disc, $h_s$, the mass of the stellar bulge, $M_{\text{bulge}}$, the initial metallicity of the gas in the disc, $Z_{\text{disc}}$, (the CGM initially contains no metals), the initial temperature of the disc, $T_{\text{disc}}$.

<table>
<thead>
<tr>
<th></th>
<th>Small</th>
<th>Fiducial</th>
<th>Large</th>
</tr>
</thead>
<tbody>
<tr>
<td>$M_{\text{tot}}$</td>
<td>$10^9 M_\odot$</td>
<td>$10^{10} M_\odot$</td>
<td>$10^{11} M_\odot$</td>
</tr>
<tr>
<td>$R_{\text{vir}}$</td>
<td>16.3 kpc</td>
<td>35.0 kpc</td>
<td>75.5 kpc</td>
</tr>
<tr>
<td>$M_{\text{disc}}$</td>
<td>$3.5 \times 10^7 M_\odot$</td>
<td>$3.5 \times 10^8 M_\odot$</td>
<td>$3.5 \times 10^9 M_\odot$</td>
</tr>
<tr>
<td>$f_{\text{gas}}$</td>
<td>0.5</td>
<td>0.5</td>
<td>0.5</td>
</tr>
<tr>
<td>$r_s$</td>
<td>0.33 kpc</td>
<td>0.70 kpc</td>
<td>1.52 kpc</td>
</tr>
<tr>
<td>$h_s$</td>
<td>33 pc</td>
<td>70 pc</td>
<td>152 pc</td>
</tr>
<tr>
<td>$M_{\text{bulge}}$</td>
<td>$3.5 \times 10^6 M_\odot$</td>
<td>$3.5 \times 10^7 M_\odot$</td>
<td>$3.5 \times 10^8 M_\odot$</td>
</tr>
<tr>
<td>$Z_{\text{disc}}$</td>
<td>0.1 $Z_\odot$</td>
<td>0.1 $Z_\odot$</td>
<td>0.1 $Z_\odot$</td>
</tr>
<tr>
<td>$T_{\text{disc}}$</td>
<td>$2.1 \times 10^3$ K</td>
<td>$10^4$ K</td>
<td>$4.6 \times 10^4$ K</td>
</tr>
</tbody>
</table>

2.3 Simulations

2.3.1 Initial conditions and simulation details

We simulate isolated galaxies comprising of a stellar and gas disc, a stellar bulge, a hot gaseous atmosphere and a static background potential representing the dark matter component. The dark matter follows an NFW profile (Navarro et al. 1997) with concentration parameter $c = 10$ and spin parameter $\lambda = 0.04$ for all galaxies simulated. The baryonic component is generated using MAKE NEW DISK (Springel et al. 2005). The disc density profile is exponential in radius. The stellar disc has a Gaussian vertical density profile with a scale height 0.1 times the scale radius. The stellar bulge has an Hernquist (1990) density profile with a scale length of 0.1 times the scale radius of the disc. The collisionless particles comprising the stellar disc and bulge in the initial conditions do not contribute to stellar feedback. The vertical structure of the gas disc is determined so as to obtain initial hydrostatic equilibrium.

We simulate three galaxies in this work with properties described in Table 2.1. The majority of simulations in this work are of a galaxy with a total mass of
Table 2.2 For the three resolutions adopted, from left to right: the target mass of cells (and star particles), the number of gas cells (excluding the CGM) and star particles in the initial conditions for the fiducial galaxy, cell diameter at the star formation threshold density of $10 \text{ cm}^{-3}$ (note that due to our Lagrangian method, cells can become much smaller, with $\Delta x \propto \rho^{-1/3}$) and minimum gravitational softening lengths for gas cells (and fixed softening length for star particles).

<table>
<thead>
<tr>
<th>Resolution</th>
<th>$m_{\text{cell}}$ [M$_{\odot}$]</th>
<th>$N_{\text{part}}$</th>
<th>$\Delta x_{\text{SF}}$ [pc]</th>
<th>$\epsilon_{\text{min}}$ [pc]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Low</td>
<td>2000</td>
<td>$1.925 \times 10^5$</td>
<td>22.7</td>
<td>8.1</td>
</tr>
<tr>
<td>Intermediate</td>
<td>200</td>
<td>$1.925 \times 10^6$</td>
<td>10.6</td>
<td>3.8</td>
</tr>
<tr>
<td>High</td>
<td>20</td>
<td>$1.925 \times 10^7$</td>
<td>4.9</td>
<td>1.8</td>
</tr>
</tbody>
</table>

$10^{10}$ M$_{\odot}$. This is analogous to a local dwarf, for example the Small Magellanic Cloud prior to falling into the Milky Way (MW), or a high redshift MW progenitor. We refer to this setup as the fiducial galaxy. This setup is comparable to the G8 model in Rosdahl et al. (2015). We also simulate two additional systems, ‘Small’ and ‘Large’, which are an order of magnitude lower and higher in mass, respectively, which can again be thought of either as local faint dwarf/intermediate-mass discs or as high redshift analogs. For our fiducial model, we initialise the disc with a temperature of $10^4$ K. We scale the initial disc temperature with the virial temperature of the halo for the ‘Small’ and ‘Large’ models. This is to avoid an initially vertically diffuse disc for the ‘Small’ model, while also maintaining consistency between the models$^5$. The gas in the disc is initialised with a metallicity of $Z = 0.1$ Z$_{\odot}$. To roughly represent the CGM we include a hot gas atmosphere of uniform density $n_H = 10^{-6}$ cm$^{-3}$, uniform temperature $10^6$ K and zero metallicity.

Gas cells and star particles (both those present in the initial conditions and newly created stars) share the same mass, 2000 M$_{\odot}$, 200 M$_{\odot}$ and 20 M$_{\odot}$ for the low, intermediate and high resolution runs, respectively. At the highest resolution, the mass of cells/star particles approaches that of the total ejecta mass per SN (10 M$_{\odot}$). We have confirmed that the refinement/derefinement scheme in place in AREPO is sufficiently effective such that star particles always have enough mass to provide the full ejecta mass budget in all but a negligible fraction of SN events ($\lesssim 1\%$). Tables 2.2 and 2.3 contain details of every simulation presented in the main body of this work, including the galaxy model used, the resolution and number

$^5$While the discs will rapidly cool once the simulation start and the vertical structure will settle into an equilibrium configuration, we find that if the initial disc structure is too vertically diffuse in the ‘Small’ model, the resulting collapse is too severe and does not allow the disc to settle satisfactorily.
Table 2.3 Details of all simulations presented in this chapter. From left to right we list: galaxy model used (see Table 2.1), feedback method used (and any additional information), feedback and star formation parameters, and total newly formed stellar mass present at 250 Myr.

<table>
<thead>
<tr>
<th>Galaxy</th>
<th>Feedback</th>
<th>Parameters</th>
<th>( M_\ast \left[ 10^7 , M_\odot \right] )</th>
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<td>( f_{\text{kin}} = 0.28 )</td>
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<tr>
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<tr>
<td>Fiducial</td>
<td>Delayed cooling</td>
<td>( f_{\text{kin}} = 0.0, \sigma_{\text{thresh}} = 10 , \text{km s}^{-1} ) ( t_{\text{diss}} = \Delta x / \sigma_{\text{FB}} ) (variable ( t_{\text{diss}} ))</td>
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</tbody>
</table>
of cells/particles, gravitational softenings, additional star formation and feedback parameters and the mass of new stars formed after 250 Myr.

2.3.2 Disc morphologies and gas phases

Fig. 2.2 shows face-on and edge-on projections of the gas and newly formed stars in the highest resolution simulations after 250 Myr. Without feedback, the gas disc cools efficiently and adopts a highly clumpy morphology on large scales. The gas in these clumps is extremely efficient at forming stars. Hence, the distribution of newly formed stars also follows this clumped morphology. There is no single dominant bulge component, instead there are multiple large clumps near the centre of the disc. Seen edge-on, the gas disc is very thin as, having cooled, it lacks vertical pressure support. The morphology for the simulations without feedback is similar to those carried out at lower resolutions (see Appendix 2.D).

The thermal, mixed (not shown) and kinetic feedback schemes are able to prevent the formation of gas clumps, instead forming complex structures of dense gas and spiral arms. This structure is also reflected in the disc of newly formed stars\(^6\). The multiple clumps of stars seen in the no feedback case are not present, though there is a definite overdensity of new stars in the centre of the disc. The global surface density of newly formed stars is greatly reduced (see Section 2.3.3). Seen edge-on, a complex vertical gas structure is evident with outflows present. At lower resolutions, this morphology is not evident (see Appendix 2.D, Figs. 2.21 and 2.22 for equivalent projections). Instead, the thermal, mixed and kinetic feedback schemes are unable to prevent the formation of dense clumps of gas. The subsequent evolution of the disc is then broadly similar to that of the runs without feedback. This clearly indicates that a mass resolution of at least \(\sim 20 \, M_\odot\) is needed for these feedback schemes to become effective.

The simulation with delayed cooling using a fixed dissipation time results in a completely disrupted disc. When the first SNe occur, they are able to eject most of the gas from the centre, leaving behind a central, low density region at 250 Myr, as evident in Fig. 2.2. The projection of newly formed stars shows an unusual ring-like structure. This is caused by the violent ejection of gas from the centre, forming stars in areas of compression as the resulting shock is transmitted through the disc plane, essentially leading to a ‘positive’ feedback. This behaviour is also apparent in the lower resolution simulations but can largely be regarded as

\(^6\)The slight ring-like structure apparent near the edge of the discs is a result of the disc settling from the initial conditions. This is a common feature of isolated disc simulations and does not affect the results.
Fig. 2.2 Projections of gas and newly formed stars at 250 Myr for different feedback runs at 20 $M_\odot$ resolution, viewed both face-on and edge-on. The mixed feedback simulation is not shown as the results are similar to the thermal and kinetic feedback simulations. The simulation without feedback results in dense clumps of gas which produce stars at a high rate. The simulations with classical, delayed cooling with variable $t_{\text{diss}}$ and mechanical feedback schemes are able to suppress the formation of dense clumps and reduce the mass of stars formed. They all show very similar disc morphologies with gas and stars exhibiting spiral patterns. The delayed cooling scheme is far too effective and blows up a large fraction of the gaseous disc leading to ring-like structures of newly formed stars. Equivalent plots for the lower resolution simulations can be found in Appendix 2.D, Figs. 2.21 and 2.22.
a numerical artifact. The strength of the feedback, and resulting gas and stellar morphologies, indicate that the choice of parameters for this scheme is not appropriate and further tuning is necessary. Further discussion of the issues with the delayed cooling in our simulation can be found later and simulations carried out with different parameters can be found in Appendix 2.C.

The delayed cooling run with variable dissipation time ($t_{\text{diss}} = \Delta x / \sigma_{\text{FB}}$) is not as strong. The disc morphology is similar to the thermal, mixed and kinetic feedback schemes at this resolution, with suppression of large scale clumping without destruction of the disc. However, perhaps counter-intuitively, this feedback scheme becomes stronger at lower resolutions (as evidenced by Figs. 2.21 and 2.22), disrupting the disc in the 2000 M$_\odot$ simulation. This is because the dependence on the cell diameter results in very short dissipation times at high resolution. At our highest resolution of 20 M$_\odot$, the cooling is essentially not delayed at all, resulting in a straight thermal dump and hence similarity to the classic thermal feedback scheme. Qualitatively, this is the desired behaviour, with the delayed cooling being reduced at resolutions high enough to resolve the Sedov-Taylor phase, but increased at lower resolution. However, the lack of convergence with resolution (and disc destruction at low resolution) indicates some form of tunable parameter might need to be introduced to refine this scheme.

Finally, the mechanical feedback scheme results in morphologies similar to the classical feedback schemes. Uniquely among the schemes tested, the mechanical feedback is able to produce these morphologies across two orders of magnitude in resolution. While the classical feedback schemes overcool at lower resolutions, the mechanical feedback scheme is still able suppress large scale clumping and the formation of high density gas, without destroying the disc. Unlike the variable dissipation time delayed cooling scheme, the implicit modulation of small scale feedback strength as a function of resolution in the mechanical scheme is able to produce convergent disc morphologies at our three resolutions.

Fig. 2.3 shows phase diagrams for the gas in the highest resolution simulations at 250 Myr (similar plots for the lower resolution runs may be found in Appendix 2.D). In the no feedback simulation, the majority of the gas in the disc has cooled well below $10^2$ K and there is a substantial quantity of gas at high density, as far as $\sim 10^6$ cm$^{-3}$. Once the gas enters this cold, dense phase, the resulting evolution is regulated by the non-thermal pressure floor and is highly dependent on the choice of parameters (for more details see Appendix 2.A), though our results are qualitatively similar to comparable simulations in the literature (e.g. Rosdahl et al. 2015, 2017; Hu et al. 2016, 2017). At lower resolution, the results are similar, though gas does
Fig. 2.3 Phase diagram for gas within the virial radius at 250 Myr for different feedback simulations at 20 M⊙ resolution. Colour coding is according to the fraction of mass in a given pixel. The vertical dotted line shows the star formation density threshold, n_{SF}. The region of the phase diagram below the diagonal dashed line is where the pressure is dominated by the non-thermal Jeans pressure floor, rather than conventional thermal pressure. The mixed feedback simulation is not shown as the results are similar to the thermal and kinetic feedback simulations. While the majority of gas resides at low temperature and high densities, i.e. within the disc, all feedback models are able to remove a fraction of gas from the ISM and to heat it to high temperatures above 10^4 K launching galaxy-scale outflow. The delayed cooling launches a large outflow at early times, the majority of which is outside the virial radius at 250 Myr. The plot for the mechanical feedback simulation is labelled, showing the location of disc (gas denser than 10^{-4} cm^{-3} is within 3 scale radii and heights), outflowing material (the region marked on the plot is all outflowing at more than 50 km s^{-1}, but still within the disc region) and the CGM on the phase diagram. Also marked is the circulation of gas around the diagram due to the galactic fountain effect.

not reach quite such high densities, an expected consequence of lower resolution. At the highest resolution, the classical feedback schemes are able to maintain a warm phase in the disc. While cold, dense gas is still present it does not reach the high densities seen in the no feedback simulation. An additional component of gas is apparent on the phase diagrams, between T ∼ 10^4 − 10^7 K and n ∼ 10^{-5} − 10 cm^{-3}. This is gas that has received feedback energy and is expanding up out of the plane of the disc. When viewed as a function of time, a cyclical pattern anticlockwise around the phase diagram may be observed with gas cooling and contracting to star...
forming densities (moving down and right), being heated by feedback (moving up), expanding (moving left). Gas which rains back on the disc (the so-called galactic fountain, see Section 2.3.3) cools and drops back down the phase diagram and may enter the cycle again.

The phase diagram for the delayed cooling with fixed dissipation time feedback at 250 Myr shows a complete absence of dense gas. In addition, because the feedback has efficiently quenched star formation (see Section 2.3.3), there are no further SNe after the initial budget has been exhausted. This results in the lack of gas above $10^4$ K (with the exception of that in the CGM). When the delayed cooling scheme is used with a variable dissipation time, at the highest resolution the phase diagrams are similar to those for the classical schemes. However, at lower resolution as the feedback becomes stronger, they become similar to the delayed cooling with fixed dissipation time. In the highest resolution simulation, the mechanical feedback scheme produces phase diagrams similar to the classical schemes, although the highest density gas has been curtailed. The phase diagrams look similar across all resolutions.

### 2.3.3 Star formation rates and outflows

Fig. 2.4 shows the mass of stars formed and the star formation rates as a function of time for all feedback schemes at all three resolutions. The no feedback simulations are similar across all three resolutions. As gas cools and reaches star forming densities, star formation begins. After a sudden jump in star formation at the beginning of the simulation, the SFR rises graduallly to a roughly constant rate $\sim 0.2 - 1$ M$_\odot$ yr$^{-1}$. The higher resolutions result in denser clump formations leading to slightly higher SFR on average, but the total stellar mass formed is similar between the three resolutions ($7.64 \times 10^7$ M$_\odot$, $9.56 \times 10^7$ M$_\odot$ and $9.23 \times 10^7$ M$_\odot$ for the $2000$ M$_\odot$, $200$ M$_\odot$, and $20$ M$_\odot$ resolutions, respectively). It should be noted that in the absence of effective feedback, the SFR and its time dependence become regulated by the choice of non-thermal pressure floor since that impacts the scale of fragmentation and the densities reached. However, our scaling of the pressure floor with resolution (as described in Section 2.2.2) results by construction in reasonably convergent behaviour with resolution, though it is not convergent with choice of pressure floor parameter (see Appendix 2.A).

In the lower resolution simulations, the classical feedback schemes are unable to suppress star formation by more than $\sim 20\%$ in the best case as they catastrophically overcool and follow the same behaviour as the no feedback simulations. A slight trend of increased effectiveness with increased $f_{\text{kin}}$ is apparent but the total
Supernova feedback in isolated simulations of galaxy formation

Fig. 2.4 Newly formed stellar mass (top) and SFRs (bottom) as a function of time for our three resolutions. At the low and intermediate resolutions, the classical feedback schemes experience the overcooling problem and barely suppress star formation relative to the no feedback simulations. However, at the highest resolution they are able to suppress star formation. The delayed cooling schemes are in general too powerful, completely quenching star formation (with the exception of the highest resolution variable $t_{\text{diss}}$ run which is barely delaying cooling in this regime). The mechanical feedback scheme suppresses star formation by a similar amount across all three resolutions, demonstrating reasonable convergence, while also being comparable to the classical schemes at the highest resolution, suggesting it is converging onto the ‘correct’ physical result.

The impact on SFRs is weak. However, at our highest resolution of $20 \, M_\odot$, the classical feedback schemes become effective, reducing SFR and total stellar mass by around an order of magnitude with respect to the no feedback simulations.

The delayed cooling scheme with a fixed dissipation time efficiently quenches star formation with a single period of SNe activity, expelling all star forming gas from the centre of the system. The results are well converged with resolution, forming almost exactly the same stellar mass. The total stellar mass formed is only a few percent of the no feedback case. As described in the previous section, the delayed cooling with variable dissipation time results in stronger feedback at lower resolution. This can be seen in Fig. 2.4 where the SFR is similar to the fixed dissipation time simulation at $2000 \, M_\odot$, higher at $200 \, M_\odot$ and close to the classical
schemes at $20\,M_\odot$.

At the lowest resolution, the mechanical feedback scheme results in a steady star formation rate below $10^{-1}\,M_\odot\,yr^{-1}$, suppressing total stars formed by approximately a factor of 5. With increasingly higher resolution, the SFRs are suppressed slightly more, but encouragingly the total stellar mass formed is within a factor of 2 from the lowest to the highest resolution simulations. The SFR exhibits a slight dip in the $200\,M_\odot$ resolution simulations, increasing slightly towards $200\,Myr$. This is due to a ‘galactic fountain’ effect, with gas launched from the disc returning and forming new stars, which we discuss in greater detail below. At the highest resolution, the mechanical feedback scheme is reasonably similar to the classical schemes because the Sedov-Taylor phase is resolved in the majority of SNe events, so the momentum boost factor is close to unity. The mechanical scheme is slightly stronger than the classical schemes due to the few SNe at this resolution that are not sufficiently resolved by the classical schemes because they occur at high densities (for further details see Section 2.3.5).

While the general features of SFRs are converged for feedback schemes that do not overcool across different resolutions, such as the ‘Mechanical’ run, interestingly the same is not true for outflows (with the exception of delayed cooling runs, which drive strong outflows at all resolutions). Fig. 2.5 shows the total gas mass within the virial radius moving at various radial velocities for our various feedback schemes at three resolutions. In the no feedback run, the behaviour is simple and is essentially the same across all resolutions. Initially, there is a rise in gas inflow as the disc collapses vertically. Gas tagged as outflowing (moving more than $5\,km\,s^{-1}$ radially outwards) is apparent despite the lack of feedback, but this is caused by motions in the disc rather than a true outflow. After $\sim 100\,Myr$, the inflow and outflow rates are approximately equal. The gas settles in the disc plane with small net motions due to the movements of the clumps and the gas reservoir is rapidly converted to stars. Note that in a cosmological context, there would be a constant net inflow from well outside the disc due to cosmic accretion. In our setup, the initially inflowing gas is simply the disc settling into an equilibrium configuration as it cools, our background uniform CGM has long cooling times, largely stays in place and does not accrete onto the disc.

The same behaviour is apparent for the classical feedback schemes at the lower two resolutions. Due to overcooling, the feedback is unable to suppress the inflowing gas, i.e. neither to stabilise the disc with a larger scale height nor to drive any appreciable outflows. At the lowest resolution, there is no gas outflowing faster than $100\,km\,s^{-1}$. There is a small fraction ($\sim 0.1\%$ of the total gas mass) moving
Fig. 2.5 Mass of gas moving at various radial velocities within the virial radius as a function of time for our different feedback schemes at all three resolutions. 

*Top:* Total gas mass within the virial radius (dashed curves), mass of gas radially outflowing and inflowing over 5 km s$^{-1}$ (solid and dotted curves, respectively).

*Bottom:* mass of gas radially outflowing at more than 50, 100 and 250 km s$^{-1}$ (solid, dotted and dashed curves, respectively). Mass of the outflowing gas is very sensitive to the resolution of simulations and only in the highest resolution runs do feedback schemes launch significant outflows.

faster than 50 km s$^{-1}$ at late times, but these are merely motions of the clumpy disc and are present in the no feedback run. At the 200 M$_{\odot}$ resolution, the feedback is able to generate some outflows faster than 50 km s$^{-1}$ ($\sim$ 1 % of the total gas mass) and a small amount moving faster than 100 km s$^{-1}$ once the density in the disc has dropped slightly due to conversion of gas to stars, reducing the overcooling effect. However, this little mass moving at relatively low velocities is not able to make it far from the disc plane and the net inflow and outflow rates are still comparable. Despite significantly suppressing star formation in the lowest resolution simulation, mechanical feedback struggles to launch outflows. It generates marginally larger
outflow rates than the classical feedback schemes, but the inflow and outflow rates match after 50 Myr. In the 200 M⊙ resolution simulation, the feedback is able to drive a much stronger outflow at early times, suppressing inflow and launching a significant mass of material faster than 50 km s⁻¹ (and a smaller amount faster than 100 km s⁻¹). However, this material is not moving fast enough to escape the galaxy, so it returns back to the disc in a galactic fountain, with the inflow rates overtaking the outflow rates at around ~ 125 Myr.

At the highest resolution, the classical and mechanical feedback schemes are relatively similar and are able to launch strong outflows. From around 50 Myr onwards, over 10⁷ M⊙ of gas is moving faster than 50 km s⁻¹, the majority of which is moving faster than 100 km s⁻¹ and a non-negligible fraction is moving faster than 250 km s⁻¹. The net outflow rates are approximately an order of magnitude larger than the inflow rates at 100 Myr. However, after this point, inflow rates rise and outflow rates drop slightly as the lower velocity gas begins to stall and flow inwards. At around 180 Myr, the inflow rates exceed the outflow rates. This galactic fountain effect can be seen in Fig. 2.6, which shows gas in a vertical slice through the system colour coded by radial velocity at several times for the mechanical feedback scheme at the higher two resolutions. In the top panel, in the 200 M⊙ simulation, an outflow is launched at 50 Myr. As time progresses, the outflowing gas moves up away from the centre of the system, but gas begins to flow inwards, segregated by velocity. By 200 Myr, the centre is dominated by returning gas while only the fastest moving gas has continued to outflow. At 250 Myr, a new outflow has just been launched, resulting in a complex, interleaved pattern of gas outflowing, inflowing and outflowing with increasing height from the disc. At the higher resolution, the initial outflow is much faster. At 100 Myr and 150 Myr, the covering fraction of outflowing gas is much larger than in the lower resolution run with all but the very central regions dominated by outflows. A complex structure of outflowing gas is apparent, with a ‘finger’ like pattern of various outflow velocities. Again, at the final snapshot at 250 Myr, the central regions contain largely inflowing gas, but far more gas continues to outflow as compared to the lower resolution.

As previously discussed, the delayed cooling schemes launch very strong outflows. As can be seen in Fig. 2.5, once again the delayed cooling with fixed dissipation time is extremely well converged across the three resolutions, drastically suppressing inflows and launching large quantities of gas at high velocities from essentially a single period of SNe activity (although, inflow rates begin to pick up at late times once SNe have been shut off due to star formation quenching). However,
Fig. 2.6 Slices through the centre of the mechanical feedback simulation at 200 $M_\odot$ (top) and 20 $M_\odot$ (bottom) resolutions showing radial velocity at various times (outflowing gas is in red, while the inflowing gas is in blue). The horizontal dotted lines show the planes at 1 and 10 kpc away from the midplane of the disc, used to examine the outflows in Fig. 2.7. Not only is the outflow stronger in the higher resolution run but the spatial structure of inflowing and outflowing gas changes as well with resolution, where the galactic fountain effect is more pronounced in the lower resolution simulation.

As previously mentioned, this feedback seems unphysically strong. The delayed cooling with variable dissipation time follows similar behaviour to the fixed dissipation time scheme, but once again converges to the classical schemes with higher resolution.

Having examined the bulk mass in outflows, it is also instructive to consider the properties of the outflow at certain heights above the disc plane. Fig. 2.7 shows the outflow rates (i.e. only considering outflowing gas, not inflowing), mass loading factors and mass-weighted average outflow velocity at 1 kpc and 10 kpc above the disc as a function of time for the highest two resolutions. We calculate the mass outflow as

$$\dot{M}_{\text{out}} = \frac{\sum_i m_i v_{\text{out},i}}{\Delta z},$$

(2.23)

where the sum is over all cells within a slice (parallel to the disc plane) of thickness $\Delta z$ centred on the target height (i.e. 1 or 10 kpc) that have a positive outflow velocity $v_{\text{out}}$ (vertically away from the disc plane, rather than radially as in the
Fig. 2.7 Mass outflow rates (top), mass loading factors (middle) and mass-weighted average outflow velocities (bottom) for our different feedback models across planes at 1 and 10 kpc from the disc midplane for our two highest resolution simulations (left and right panels). The dashed grey lines indicate the escape velocity at the relative disc height. The 2000 M$_\odot$ simulations are not shown as outflows for all except delayed cooling are negligible (see Figs. 2.5 and 2.23). Outflow velocities comparable to the escape velocity and mass loading factors of a few are only reached at the highest resolution simulations for all feedback runs (except for delayed cooling runs which are over-efficient).

We adopt $\Delta z = 200$ pc. The mass loading factor, $\beta_v$ is the ratio of the mass outflow rate to the star formation rate and is essentially a measure of the efficiency of stellar feedback to drive outflows. There is of course a delay between the formation of a given stellar population and the outflow its feedback eventually drives reaching a given height. However, here we use the instantaneous ratio between $\dot{M}_{\text{out}}$ and the SFR rather than a more complex binning scheme as our SFRs are, on the whole, steady over long periods (with the exception of the delayed cooling schemes, for which the mass loading must be interpreted with some caution). Finally, we plot the mass-weighted mean outflow velocity (i.e. ignoring inflowing gas) alongside the escape velocity at that height$^7$.

$^7$We calculate the escape velocity at the relevant height directly above the centre of the disc
As demonstrated in Fig. 2.5, with 200 $M_\odot$ resolution, the classical feedback schemes are unable to drive much of an outflow. After some initial gas flow over 1 kpc as the system settles from the initial conditions (also apparent in the no feedback simulations; this is a ‘spurious’ mass loading amplified by low SFRs), the mass outflow drops off. There is a small increase after 100 Myr as the feedback becomes more efficient and the gas reservoir is used up, but very little of this outflow reaches 10 kpc. At 1 kpc, this outflow has a mass loading factor below 0.1 i.e. significantly lower than the $\beta_v \gtrsim 1$ required by observations and models (for a more detailed discussion see introduction). The mean outflow velocities (admittedly dominated by the slower moving gas) are well below the escape velocity at 1 kpc. The mixed feedback has a higher mean outflow velocity at 10 kpc, however its seemingly increased effectiveness over the other methods can be put down to stochasticity amplified by the exceedingly small mass of gas that is actually outflowing at that distance. The mechanical feedback is able to generate a slightly more vigorous outflow, with a mass loading factor between $1-10$ at 1 kpc. However, it is unable to sustain the outflow as previously discussed, with gas returning in a galactic fountain. Again, very little of the outflow reaches 10 kpc.

At the higher resolution of 20 $M_\odot$, the classical and mechanical feedback schemes are able to launch much stronger, sustained outflows. Once the outflow has reached the heights we are investigating, mass loading is around 10 at 1 kpc and over unity at 10 kpc. The mean velocities are below the escape velocity at 1 kpc, but a significant quantity of gas (see Fig. 2.5) is moving much faster. By 10 kpc, the slower moving gas having begun to drop back to the disc, the mean velocities are comparable to the escape velocity.

The delayed cooling simulations with fixed dissipation time are able to launch strong, but short lived outflows. Having completely quenched star formation, there is no source for additional driving of outflows beyond the initial burst. The instantaneous mass loading factor becomes an unreliable metric in such conditions, since it must naturally tend to infinity as SFR tends to zero, however we plot it here for reference. Again, using the variable dissipation time results in strong outflows at lower resolutions but similar results to the classical and mechanical feedback schemes at the highest resolution.

Fig. 2.8 contains vertical slices through the disc at 250 Myr for the highest resolution simulations, showing gas density, temperature and metallicity. Generating no outflows, the no feedback simulation shows a cold, thin, dense disc. The from the initial conditions. Deviations from the initial conditions during the course of the simulation have a negligible impact on $v_{\text{esc}}$ at 1 kpc and are insignificant at 10 kpc as the dominant component is the static halo potential.
Fig. 2.8 Density, temperature and metallicity slices at 250 Myr for 20 $M_\odot$ resolution runs. The no feedback simulation unsurprisingly launches no outflows, leaving a high concentration of metals in the centre of the system. The classical, delayed cooling with variable $t_{\text{diss}}$ and mechanical schemes all produce comparable outflows (differences are largely due to variability with time). The outflows have a complex filamentary structure and are multiphase with temperatures spanning $\sim 6$ orders of magnitude. The delayed cooling scheme evacuates gas from the centre of the system, leaving a column of very hot, low density gas. Equivalent plots for the lower resolution simulations, which are unable to launch such strong outflows, can be found in Appendix 2.D, Fig. 2.23.

central regions have very high metallicities since the ejecta from SNe stay within the star forming regions. The thermal, mixed (not shown), kinetic, mechanical and variable dissipation time delayed cooling feedback schemes have qualitatively similar outflows. The outflows are multiphase (with temperatures in the range $\sim 10^2 - 10^7$ K) and have a complex structure, with many individual filaments apparent. Observations of galactic outflows reveal them to be comprised of mul-
tiphase gas: molecular gas at $\sim 10 - 10^3$ K observed at radio wavelengths (e.g. Walter et al. 2002, 2017; Bolatto et al. 2013), material around $\sim 10^4$ K observed in the optical and near-UV (e.g. Pettini et al. 2001; Martin et al. 2012; Soto et al. 2012) and $\sim 10^7 - 10^8$ K plasma seen in X-rays (e.g. Martin 1999; Strickland & Heckman 2007, 2009). At approximately the peak of the outflow (150 Myr), for the mechanical feedback simulation, in material moving radially outwards at more than $100$ km s$^{-1}$, the proportions of cold ($< 2000$ K), warm ($2000 - 4 \times 10^5$ K) and hot ($> 4 \times 10^5$ K) material are 3.1%, 78.9% and 17.9% by mass, respectively, or 0.1%, 56.8% and 43.1% by volume. Thus, while the dominant fast moving wind component is warm, a cold component is present in the outflow. The cold component dominates by mass in material moving outwards between $5 - 100$ km s$^{-1}$, with the proportions of cold, warm and hot components being 56.4%, 39% and 4.6% respectively (by volume, 0.6%, 4.2% and 95.2%, the very slow moving CGM material dominating the hot component here). As the outflow progresses and the galactic fountain effect becomes apparent, the proportions remain similar. The cold component dominates the returning gas, but warm gas also returns. It appears that the returning cold component contains both initially cold outflowing gas as well as material from the warm phase that has cooled. In summary, we find that the cold gas mainly traces the lower velocity outflows while the warm and hot medium probe the faster moving outflows; if this effect is present in real galaxies, observing one component alone will give a biased measurement of the outflow velocity.

The results are similar for the thermal, mixed, kinetic and variable dissipation time delayed cooling feedback schemes at this resolution. Small differences between schemes apparent in the figure are largely due to stochasticity. The outflowing gas is enriched with metals and there is a dependence of metallicity on opening angle. The most metal enriched regions of the outflow are in the centre ($Z \gtrsim 0.25Z_\odot$), containing the highest concentration of SNe ejecta, whereas towards the edges of the outflow the metallicity is closer to the initial disc gas metallicity of $0.1Z_\odot$. The kinetic feedback simulation shows a high metallicity outflow of $\sim 10^6$ K gas, having had an outflow event shortly before 250 Myr. The delayed cooling with fixed dissipation time simulation exhibits a column of low density, high temperature gas that extends all the way though the centre of the disc, with cold, denser gas building up on the fringes of the outflow. At lower resolutions (see Fig. 2.23), the outflows are much weaker (as described above) for all simulations except the delayed cooling schemes. At 200 $M_\odot$ resolution, the classical feedback schemes launch warm $10^4$ K outflows with dense, cold edges at the interface with the CGM. The outflows are highly enriched because the mass loading is so low i.e. a substantial amount of SNe
occurred to launch the outflows.

2.3.4 The Kennicutt-Schmidt relation

The link between gas surface density, $\Sigma_{\text{gas}}$, and SFR surface density, $\Sigma_{\text{SFR}}$, is an important diagnostic of star formation in galaxies. Specifically, the Kennicutt-Schmidt relation, $\Sigma_{\text{SFR}} \propto \Sigma_{\text{gas}}^{1.4}$ (Kennicutt 1998), has been well established by observations of galaxies in the local Universe. Thus, in addition to suppressing absolute SFRs, it is necessary for simulations to simultaneously reproduce this relation. It is possible to have very different values of $\Sigma_{\text{SFR}}$ for the same global $\Sigma_{\text{gas}}$, dependent on the small scale star formation and the degree of clustering in star formation. Our choice of small scale star formation law to some extent impacts the resulting global KS relation. For example the choice of $\dot{\rho}_* \propto \rho/t_{\text{ff}} \propto \rho^{3/2}$ generally leads to the correct slope, but this does not guarantee the correct normalization, as shown below.

Fig. 2.9 shows the global star formation rate surface density, $\Sigma_{\text{SFR}}$, as a function of global gas surface density, $\Sigma_{\text{gas}}$ for our simulations, each point representing one of the simulations at a particular time (points are evenly spaced by 25 Myr between 25 and 250 Myr). We define the surface densities as

$$\Sigma_{\text{SFR}} = \frac{\dot{M}_* (< R_{\text{SFR},90\%})}{\pi R_{\text{SFR},90\%}^2},$$

(2.24)

and

$$\Sigma_{\text{gas}} = \frac{M_{\text{gas}} (< R_{\text{SFR},90\%})}{\pi R_{\text{SFR},90\%}^2},$$

(2.25)

where $R_{\text{SFR},90\%}$ is the disc radius enclosing 90% of the total SFR.8

For comparison, we plot global measurements from 61 normal spirals (Kennicutt 1998), similar global measurements from 19 low surface brightness galaxies (Wyder et al. 2009) and sub-kpc observations of 18 nearby galaxies (Bigiel et al. 2008). For reference, we also plot the power law fit with a slope of 1.4 from Kennicutt (1998), however it is worth noting that the slope is possibly too shallow for this range of measurements. This fit was made simultaneously to the 61 spirals plotted as well as 36 higher surface density starburst galaxies (not plotted). At lower surface densities, the relation appears to steepen, possibly due to some form of star

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8Results are fairly insensitive to the choice of the fraction of the SFR enclosed, merely sliding points up and down the Kennicutt-Schmidt relation. We only include gas within 2 kpc of the disc plane, although our results are insensitive to removing this constraint because the gas surface density is completely dominated by mass near the disc plane.
Fig. 2.9 SFR surface density plotted as a function of gas surface density for different feedback runs at all three resolutions (as indicated by different symbols). Each symbol represents the entire galaxy at one time between 25 - 250 Myr, the open symbols corresponding to the final snapshot at 250 Myr. The black crosses are global measurements of normal spirals from Kennicutt (1998), while red crosses are similar measurements for low surface brightness galaxies from Wyder et al. (2009). The contour is derived from multiple sub-kpc measurements of 18 galaxies, including spirals and dwarfs from Bigiel et al. (2008). We plot here the contour corresponding to more than 5 data points per 0.05 dex-wide cell. For reference, we also plot the power law with a slope of 1.4 from Kennicutt (1998), fitted to both the data plotted here and higher surface density starburst galaxies. While classical feedback schemes agree with the observed Kennicutt-Schmidt relation only at the highest resolution, the mechanical feedback produces realistic SFR and gas surface densities at all three resolutions.

formation threshold (e.g. Kennicutt 1989; Martin & Kennicutt 2001; Bigiel et al. 2008). Thus, it makes more sense to compare our results to the data points rather than the fit plotted.

Except at the highest resolution, the no feedback and classical feedback simulations all lie well above the observed relation, although once the system has finished clumping after $\sim 100$ Myr, the points have approximately the correct slope. This is mainly due to the small scale star formation law adopted forcing $\dot{\rho}_* \propto \rho^{3/2}$. The simulations then progress to lower SFR and gas surface densities as the gas
reservoir is consumed. At the highest resolution, the classical schemes are able to quench star formation efficiently and so drop into agreement with observations. Relative to the classical feedback schemes, the other three feedback mechanisms produce an order of magnitude lower SFR surface densities for the same gas surface density, lying close to the observed relation at all three resolutions. The delayed cooling with fixed dissipation time efficiently destroys the disc, so the majority of the snapshots lie outside the range of the plot. The same is true of the variable dissipation time run at the lower resolution, though the high resolution run is well within the observed points. The mechanical feedback runs at all resolutions agree well with the observations. The simulations track up and down the relation with time as the gas surface density changes, partly due to gas consumption, but mainly due to outflows. For example, the cluster of $20\, M_\odot$ resolution mechanical feedback points (yellow squares) near the bottom left of the relation correspond to the period after the peak of the outflow at about 100 - 200 Myr, but returning gas from the galactic fountain causes the disc to have moved back up the relation by 250 Myr (open yellow square). In addition to variation over time, this effect also causes the differences between the three resolutions: the lower resolutions tend to lie higher up the relation because their weaker outflows do not drop the disc surface density as much. Despite this, the mechanical feedback points all lie close to the observations even though their exact position on the relation varies with resolution, this difference caused by the failure of resolution convergence with respect to outflows.

### 2.3.5 Host sites of star formation and supernovae

Fig. 2.10 shows PDFs of the local densities where stars are formed (top panels) and SNe explode (bottom panels) for our different feedback runs and at all three resolutions. Looking at the sites of star formation in the simulations without feedback, a double peak form is apparent (though in the lowest resolution, the lower density peak is suppressed into more of a tail). This shape is a consequence of using a star formation threshold density (indicated by a vertical dashed line in Fig. 2.10). At the beginning of the simulation, as gas densities increase and cross this threshold, the first burst of star formation occurs, building a peak in the PDF just above the threshold density. The gas continues to clump until it reaches the maximum density the resolution and pressure floor allows. The majority of star formation then occurs at this density, building a high density peak in the PDF, enhanced by the fact that the SFR is higher in denser regions (i.e. $\dot{\rho}_* \propto \rho^{3/2}$). The sites where SNe occur (in this case, where mass and metals are returned but no feedback energy is deposited) are therefore an almost direct mapping from the
Fig. 2.10 PDFs of the densities of the sites where stars are formed (top) and where SNe occur (bottom) throughout the entire simulation. In addition, the star formation density threshold is marked with a vertical dashed line. Without efficient feedback, the majority of stars form at high densities and SNe occur in these regions. If the feedback is able to disrupt the dense birth clouds, then subsequent SNe occur at much lower densities, leading to a tail in the PDF well below the star formation density threshold. This also prevents star forming gas from reaching such high densities.

star formation PDF because the local ISM is essentially unchanged from the star particle being born to its SNe events occurring (although continued star formation will act to drop the local density by transferring gas mass into star particles, while gas may continue to collapse to higher densities before the first SNe occur).

Inefficient feedback (i.e. the classical schemes at the lower two resolutions) follow the same behaviour as they are unable to disrupt the dense clumps of gas where the star particles are formed. In contrast, efficient feedback is able to prevent the gas from clumping to high densities, therefore increasing the fractional contribution of lower density star formation. In the case of very strong feedback that disrupts the disc (i.e. delayed cooling) the PDF of star formation is entirely dominated by the first burst of star formation, before the SNe that those star particles produce completely quench star formation. For more moderate feedback runs, the effect of a cycle of varying SFRs, caused by the return of previously ejected mass (the
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galactic fountain), can allow a building up of a small peak at high density. For example, in the 200 \( M_\odot \) resolution mechanical feedback run, after a period of low SFR between 50 and 200 Myr (see Fig. 2.4), the small peak at \( \sim 10^4 \text{ cm}^{-3} \) is able to form before the feedback is able to start destroying clumps again. The degree to which this effect occurs is an indication of how effective a given feedback scheme is at dispersing dense gas at a low local SNe rate, since this is directly linked to the local SFR (with an offset arising from a delay between massive star formation and SNe occurrence). In other words, once the SFR has been reduced by feedback, if the resulting lower SNe rate is unable to efficiently prevent the return of gas, a high density peak in the SF PDF will occur. Specifically, in the highest resolution simulations, the thermal feedback shows a small peak at \( \sim 10^4 \text{ cm}^{-3} \) while the mechanical feedback does not (with mixed, kinetic and delayed cooling with variable dissipation time lying between the two in order of their effectiveness).

Note that the shape of the PDF above the star formation threshold will also be dependent upon the details of the small scale star formation prescription adopted; this is discussed in Section 2.3.8.

With efficient feedback, the shape of the PDF of densities of SNe sites above the star formation threshold corresponds closely to the star formation PDFs, as in the inefficient feedback case. However, the PDF extends well below the star formation density threshold. Since by definition the star particles from which the SNe are occurring cannot have been formed at these densities, these SNe are occurring after previous SNe have disrupted the star forming regions of their birth cloud. Of course, more efficient feedback results in more SNe occurring in low density environments. These subsequent SNe are themselves more efficient because momentum input into the ISM is higher at lower ambient densities (but also numerically, in the case of the classical schemes, because it is easier to resolve the Sedov-Taylor phase).

Thus, particularly when other faster acting stellar feedback effects are not included (as in our work), a major requirement of an efficient SN feedback scheme is that the first SNe to occur in a star forming region are able to disperse the dense gas to allow the efficiency of later SNe to be increased. The classical feedback schemes are only able to achieve this at the highest resolution probed. The mechanical feedback scheme is more successful at all resolutions, but with increasing resolution more SNe go off in lower density environments. At the highest resolution, the shape

An alternative mechanism by which SNe can occur outside of dense star forming regions requires the SNe progenitors to have moved out of their birth clouds (i.e. OB runaways, see e.g. Conroy & Kratter 2012), most likely as a result of interactions with other stars. Because we do not resolve the dynamics of individual stars in their clusters, this effect is not present in our simulations. We could adopt an additional sub-grid recipe to replicate this (e.g. Ceverino & Klypin 2009; Kimm & Cen 2014; Kim & Ostriker 2017), but this is beyond the scope of this work.
of the PDF below the star formation density threshold is similar for all feedback schemes. Because the delayed cooling schemes result in a rapid clearing of gas from the centre of the system, most SNe occur in low density gas at all resolutions.

### 2.3.6 Structure and kinematics

Fig. 2.11 shows the spherically averaged radial profiles of number density and mass fraction of gas together with new stars, and of gas separately, as well as mass-weighted gas temperature and metallicity, for the highest resolution simulations at 250 Myr. The profiles from the initial conditions are also plotted for comparison. The no feedback shows an enhancement in baryon number density and mass fraction, due to a large centrally positioned clump (see Fig. 2.2). A smaller peak caused by another clump further out is also apparent. However, the density and mass fraction of the gas component taken on its own are significantly reduced from the initial conditions, indicative of a major conversion of gas to stars in situ. The profiles are unchanged at large radii as there has been no outflow of material. The temperature within the disc gas has dropped by several orders of magnitude due to metal cooling. The metallicity of disc gas has increased by a factor of \(10^{-100}\) because of the high SFR in conjunction with the lack of outflows, resulting in very short cooling times.

The classical, mechanical and delayed cooling with variable dissipation time feedback schemes are similar to the initial conditions with respect to the baryon number density and mass fraction in the central regions. The gas mass fraction has been reduced, as have the central densities, partly due to conversion of gas to stars (as in the no feedback case) but mainly due to outflowing material. This outflow material can be seen at larger radii, where the gas mass fraction outside \(\sim 4\) kpc has been significantly increased. The gas within the disc has been prevented from runaway cooling and the average temperature is mostly between \(10^3 - 10^4\) K. Variations in temperature between the different feedback schemes are largely transient and stochastic, particularly at small radii (where the average is over less gas mass). Temperatures are significantly reduced from the initial conditions outside a few kpc, as colder outflowing gas displaces the hot CGM. Central metallicities are increased from the initial conditions by a factor of \(\sim 3 - 8\) (with the exception of a metallicity spike from the recent feedback event in the thermal simulation which has not yet dispersed). Metals have been transported into the region initially occupied by the CGM. As previously described, the delayed cooling with fixed dissipation time evacuates gas from the central regions extremely efficiently, resulting in a large drop in the central gas density and mass fraction, a spike in tempera-
Fig. 2.11 Spherically averaged radial profiles of number density of gas and newly formed stars, gas number density, gas temperature (top panels) and mass fraction of gas and newly formed stars, gas mass fraction and gas metallicity (bottom panels) for our different feedback runs at 250 Myr for simulations with 20 M⊙ resolution. The profiles from the initial conditions are shown with gray dotted curves.

ture of the remaining gas but a very small increase in metallicity (because there have been very few SNe). Despite the explosive nature of delayed cooling feedback, the mass-averaged temperature at the outer radii (1 – 10 kpc) is still much lower than the original CGM temperature. The lower resolutions simulations (not shown) have very similar profiles to those described above, dependent on whether how effective the feedback is (i.e. the classical schemes overcool so adopt the same behaviour as the no feedback case). We have also examined surface density profiles (not shown) and find that the results are similar to the radial profiles. Simulations
Fig. 2.12 Circular velocity profiles (left) and circularities, $\epsilon = j_z/j_{\text{circ}}$, (right) for newly formed stars at 250 Myr for our different feedback simulations at 20 $M_\odot$ resolution. We also plot the initial circular velocity profile and distribution of circularities for stars present in the initial conditions and of gas (within 3 scale radii and 3 scale heights) (right). The circular velocity profiles are largely unchanged from the initial conditions (due to the largely unchanged initial stellar disc and the static halo potential). The no feedback simulation is peaked close to the centre due to presence of a clump of gas and stars. The simulations with feedback reduce the circular velocity slightly by transporting mass outwards. All simulations have circularity distributions centred at $\sim 1$, indicative of a disc. There is no signature of a bulge component. The no feedback simulation has a broader distribution of circularities due to the highly clumped disc structure.

with feedback preserve the initial exponential density profile of the disc in terms of total baryons. The gas profiles are centrally cored relative to the initial profile and densities are enhanced at outer radii due to outflows.

Fig. 2.12 shows the circular velocity profiles and distribution of circularity parameter, $\epsilon = j_z/j_{\text{circ}}$, for newly formed stars for our different feedback runs at the highest resolution. The circular velocities are generally similar to initial conditions, because the distribution of the initial stellar disc and (small) bulge of old stars is largely unchanged (these components making up over half of the initial baryonic mass), while at larger radii the circular velocity is dominated by the static halo potential. The no feedback case shows a peak at small radii due to the centrally positioned clump remarked upon earlier. Note, as described above, this clump cannot be taken to be indicative of bulge formation. As seen in Fig. 2.2, there are multiple clumps present in the disc. The position of the clump is somewhat stochastic. We have found other simulations at a variety of resolutions to produce the same degree
of clumping, yet not necessarily with a clump positioned close enough to the centre to produce a central peak in the circular velocity profile. However, as seen in Fig. 2.12, efficient feedback results in a reduction in circular velocity, particularly at small radii, due to the transport of gas mass out to further radii.

The circularity parameter gives an indication of the degree of rotational support in a system by comparing the specific angular momentum in the $z$-direction (i.e. out of the disc plane) to that required for a circular orbit at the same radius. Thus, stars belonging to a disc will have $\epsilon \sim 1$, whereas a non-rotating spheroid would have a symmetric distribution about $\epsilon = 0$. All simulations have a peak around $\epsilon \sim 1$, indicative that the newly formed stars form a disc, while there is no clear signature of a bulge component (such as the small enhancement at $\epsilon = 0$ present in the initial stars). This is not surprising because the stars have formed directly from the gas disc. The distribution is peaked very slightly below $\epsilon = 1$; this is inherited from the initial gas distribution which has some degree of pressure support in addition to the rotational support (as can be seen by comparing the initial distribution of gas circularities to those of the stellar disc). The difference between no feedback and feedback simulations is apparent in the width of the distribution. With feedback, the stellar circularities are relatively narrowly distributed around 1 (irrespective of the feedback scheme adopted), whereas in the no feedback gas the distribution is considerably broader. This is caused by the highly clumpy disc formed in this run, with stars acquiring significant offsets from the circular velocity due to local interactions with clumps.

We conclude this section by remarking that a proper study into the effects of the feedback schemes on galaxy structure and kinematics should make use of galaxies formed self-consistently in a cosmological context, rather than in a disc set up ‘by hand’ as in this work. However, we find it informative to examine the extent to which an ideal system is maintained in the presence of our feedback schemes, ranging from a strongly clumped distribution in the no feedback case to total disc destruction due to overstrong feedback in the delayed cooling case.

**2.3.7 Varying galaxy mass**

In addition to our fiducial $10^{10} \, M_\odot$ galaxy, we have also run simulations of smaller ($10^9 \, M_\odot$) and larger ($10^{11} \, M_\odot$) systems without feedback and with mechanical feedback at our intermediate resolution of $200 \, M_\odot$. The results of these simulations are summarised in Fig. 2.13 alongside the equivalent simulations of our fiducial system, showing mass of newly formed stars (expressed as a fraction of the initial disc gas budget), specific star formation rates (SSFR) and mass loading factors
Fig. 2.13 A comparison of simulations with no feedback and with mechanical feedback for our three galaxy masses (see Table 2.1) at 200 $M_\odot$ resolution. Top left: Newly formed stellar mass expressed as a fraction of the initial gas disc mass. Bottom left: specific star formation rates ($\dot{M}_*/M_*$ where $M_*$ includes both old and new stellar mass). Right: mass loading factor across two planes at different distances from the disc midplane. For the fiducial galaxy, these are 1 and 10 kpc as in Fig. 2.7. For the small and large galaxies, the planes are at the same distance relative to the virial radius as in the fiducial case (0.47 and 4.7, 2.16 and 21.6 kpc, respectively). Global star formation efficiency increases with increasing system mass (a trend in line with abundance matching), though this appears to be independent of feedback in our setup. Feedback suppresses star formation by a similar factor in all three systems. Outflows become weaker with increasing system mass.
this is 1 kpc and 10 kpc, so we use 0.47 kpc and 4.7 kpc for the smaller system\textsuperscript{10} and 2.16 kpc and 21.6 kpc for the larger system.

In the smaller system, the simulation without feedback quickly starts forming stars from the beginning of the simulation, though it quickly establishes a steady SSFR of $\sim 2 \times 10^{-10} \text{yr}^{-1}$ as the smaller surface densities prevent gas from clumping to high densities. This SSFR is approximately an order of magnitude smaller than the fiducial simulation. The mechanical feedback initially follows the same evolution as the no feedback case. However, once the SFR has reached its peak, the SNe are able to unbind the gas from system due to the shallow potential well, quenching star formation. At 250 Myr, the ratio of newly formed stellar mass to the initial gas disc mass is approximately an order of magnitude lower than the fiducial simulation. This trend agrees with the general results from abundance matching that lower mass galaxies are less efficient at forming stars (below $\sim 10^{12} \text{M}_\odot$) (see e.g. Moster et al. 2013; Behroozi et al. 2013). However, the factor by which feedback has suppressed star formation relative to the no feedback simulation is similar to the fiducial case. This suggests that the lower star formation efficiency relative to the fiducial simulation is inherent to our particular setup, rather than being caused by more efficient feedback as is commonly posited to explain the phenomenon. After the spuriously high mass loading due to the low SFR at the beginning of the simulation has reduced (note that the no feedback simulation continues to have a relatively high apparent mass loading for the same reason), the mass loading at 0.47 kpc and 4.7 kpc rises dramatically as the gas is expelled from the system. The mass loading then tends to infinity because the SFR has dropped to zero (the instantaneous mass loading factor is not a good metric in a highly bursty regime).

Without feedback, the larger system follows an evolution similar to that of the fiducial system. However, it forms stars more than proportionally faster than the fiducial case, resulting in $\sim 3$ times more stellar mass formed than a simple scaling with the system mass. With mechanical feedback, the result is similar, relatively speaking, with star formation suppressed by a similar factor. Once again, the trend of a greater star formation efficiency seems to be qualitatively in line with abundance matching, but, as with the low mass system, this effect seems to be inherent to the set-up rather than caused by less efficient feedback. The mass loading factor is mostly within a factor of a few of the fiducial simulation at both distances, though is always lower. The mass loading factor is mostly below unity at 2.16 kpc and reaches a maximum of only 0.1 at 21.6 kpc, i.e. the outflows are

\textsuperscript{10}We use a thickness $\Delta z = 100 \text{pc}$ for determining the mass outflow rates for the smaller mass system.
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Fig. 2.14 Simulations with no feedback and mechanical feedback with varying star formation criteria at 200 $M_\odot$ resolution for our fiducial galaxy. We compare our fiducial values for star formation ($\epsilon_{\text{SF}} = 1.5\%$, $n_{\text{SF}} = 10$ cm$^{-3}$) with an increased star formation efficiency ($\epsilon_{\text{SF}} = 15\%$) or an increased star formation threshold density ($n_{\text{SF}} = 100$ cm$^{-3}$). *Top left:* Newly formed stellar mass. *Bottom left:* SFRs. *Right:* mass loading factor across two planes at different distances from the disc midplane. These are at 1 kpc and 10 kpc as in Fig. 2.7. Increasing $\epsilon_{\text{SF}}$ results in faster star formation, leading to much a much stronger burst of feedback which quenches subsequent star formation. Increasing $n_{\text{SF}}$ results in a similar evolution of stellar mass to the fiducial case, but produces weaker outflows.

even more inefficient than in our fiducial galaxy model, which is not unsurprising given the deeper potential well outflows need to overcome.

2.3.8 Varying star formation law parameters

While the focus of this work is on the difference between different feedback schemes, we also briefly examine here the effect of varying the parameters used with our adopted star formation law (see equation (2.2)). We rerun our 200 $M_\odot$ resolution
simulations of the fiducial galaxy without feedback and with mechanical feedback but with an increased star formation efficiency parameter, $\epsilon_{\text{SF}} = 15\%$ rather than the fiducial $1.5\%$ and also with an order of magnitude higher density threshold, $n_{\text{SF}} = 100 \, \text{cm}^{-3}$. The results are summarised in Fig. 2.14 alongside the fiducial simulations.

Increasing the star formation efficiency parameter by an order of magnitude results in the initial SFR being an order of magnitude higher than the fiducial case both with and without feedback. The simulation without feedback maintains this high SFR ($\sim 1 \, M_\odot \, \text{yr}^{-1}$), dipping slightly below the fiducial simulation’s SFR, which has risen to this value, at $\sim 110 \, \text{Myr}$ as the gas reservoir is consumed. The final newly formed stellar mass is approximately 1.25 times larger than the fiducial run. In the mechanical feedback case, the high SFR leads to a burst of strong feedback at $20 \, \text{Myr}$ which expels the gas from the centre of the system and quenches star formation. The mass in newly formed stars at $250 \, \text{Myr}$ is similar to the fiducial simulation, but the majority of stars have formed in the first $50 \, \text{Myr}$. Once again, the instantaneous mass loading factor for a regime that quenches star formation is an unreliable metric (as it tends to infinity as the SFR plummets). However, there is a brief period between $20 - 100 \, \text{Myr}$ where the SFR is non-zero, leading to a mass loading factor between $1 - 50$ at $1 \, \text{kpc}$. The outflow also easily reaches the $10 \, \text{kpc}$ plane as can be seen by the high mass loading factor.

Increasing the star formation density threshold by an order of magnitude results in an initially lower SFR in the simulation without feedback as it takes slightly longer for the gas to reach the higher star forming densities. However, by $70 \, \text{Myr}$, the SFR has reached the levels of the fiducial simulation and subsequent evolution is similar, resulting in the almost the same mass in new stars at $250 \, \text{Myr}$. The simulation with mechanical feedback is similar until $50 \, \text{Myr}$, when the SNe are able to halt further rising of the SFR. A stable SFR is established, a factor of a few higher than the fiducial simulation. The stellar mass at $250 \, \text{Myr}$ is only $1.4$ times that of the fiducial simulation. The outflow is weaker at $1 \, \text{kpc}$ by a factor of a few than the fiducial case for most of the simulation, stable at around unity. However, at $10 \, \text{kpc}$ the outflow is very weak, with mass loading factor between $10^{-3} - 10^{-2}$ (an order of magnitude lower than the fiducial simulation).

Fig. 2.15 shows PDFs of the gas densities at the sites of star formation and SNe explosions, comparing the models with altered star formation law parameters to the fiducial simulation. The simulation with increased star formation efficiency produces most of its stars in gas that is several orders of magnitude less dense than the fiducial case, because gas is rapidly converted to stars before it can collapse to
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Fig. 2.15 PDFs of the densities of the sites where stars are formed (top) and where SNe occur (bottom) for the simulations in Fig. 2.14, comparing the effects of changing the star formation law parameters using our fiducial galaxy mass and 200 M$_\odot$ resolution. Increasing $\epsilon_{\text{SF}}$ results in gas being rapidly converted to stars before it can reach high densities. A strong burst of initial feedback disrupts star forming clouds, so the majority of SNe occur at very low densities. Increasing $n_{\text{SF}}$ has very little impact on the PDFs, though slightly more SNe occur at high densities.

In the simulations with a higher star formation threshold density, without feedback, the PDFs of star formation and SNe site densities are similar, because most star formation occurs in gas at $n = 10^4$ cm$^{-3}$, well above the thresholds. However, in the runs with mechanical feedback, the situation is different. In the fiducial simulation, feedback shifts the peak star formation density down to the threshold.
In the simulation with a higher threshold density, stars are born at much higher densities. The result is that SNe occur in gas of higher density which reduces their momentum input to the ISM. The feedback is still strong enough to disrupt star forming regions, which is why the SFR is close to the fiducial case, but the reduction in momentum results in weaker outflows. Note that the reduction in momentum input is not due to overcooling, but is physical (see equation (2.22)), although the subsequent development of outflows is subject to resolution effects, as discussed in Section 2.3.3.

What we have demonstrated in this section is that changing the star formation prescription (to not unreasonable values) can have a non-negligible effect on SFR and outflows, although a more comprehensive study of these effects is beyond the scope of this work. We note that a low mass system such as our fiducial model is likely to be less robust to changes in the star formation prescription because it is easy to unbind gas if the feedback increases in strength, leaving little margin for self-regulation, although we find our results to be broadly in agreement with similar tests in Rosdahl et al. (2017). It is also likely that the inclusion of other feedback processes that act in the time between star formation and the resulting SNe might possibly help mitigate the dependence on the star formation prescription by preventing further collapse of gas (see Hopkins et al. (2011, 2013) for examples of self-regulating systems that are somewhat robust to the star formation prescription in terms of the global SFR).

2.4 Discussion

2.4.1 Comparison of SN feedback implementations

We find that our ‘classical’ schemes (simple injection of thermal energy, kinetic energy or a mixture of the two) all give very similar results. There is a slight trend for an injection of kinetic energy to result in stronger feedback but the effect is minor. At all but the highest resolution of $20 \, M_\odot$ these schemes suffer from the overcooling problem, barely suppressing star formation relative to the no feedback case and producing similar clumpy morphologies. This is not unexpected.

One can alleviate the problem slightly by injecting the energy of several SNe at once, but only up to a point. For example, Rosdahl et al. (2017) inject the energy of 40 SNe simultaneously which allows a thermal dump to be efficient in their $10^{11} \, M_\odot$ system but not in their $10^{12} \, M_\odot$ system, due to a combination of the deeper potential, stronger metal line cooling, higher densities and lower...
resolution. In addition, such an approach requires the adoption of an artificial delay time between the birth of a star particle and the triggering of a SN event. In cosmological zoom-in simulations, Aumer et al. (2013) find that their results are sensitive to this choice, particularly in the absence of other sources of stellar feedback, requiring an unphysically early value of 3 Myr to recover realistic star formation histories. With high resolution simulations of individual and clustered SNe, Kim et al. (2017) demonstrate that the evolution of a succession of SNe can differ significantly from the evolution of a single event of equivalent energy and may overestimate momentum injection. Hopkins et al. (2018) find that grouping SN events together rather than resolving them individually can lead to extremely high metallicities in dwarfs. This is because a large metal mass is injected at once which is ejected from the galaxy but is sufficiently metal-rich that it preferentially forms the next stellar population as it re-cools. This phenomenon does not occur with physically time-spaced SNe. Kimm et al. (2015) find that allowing for a realistic delay time, with individual SNe distributed between $\sim 3 - 40$ Myr, prevents the build up of dense gas prior to SNe occurring relative to a fixed delay time of 10 Myr, while also allowing later SNe to explode in low density environments produced by earlier events. The caveat, of course, is that individual SNe are more susceptible to overcooling.

Our trial of delayed cooling schemes is unsatisfactory. Unlike the other schemes explored, these schemes have adjustable parameters, which is something we wish to avoid if possible. In addition, delayed cooling schemes circumvent unphysical results caused by lack of resolution by enforcing an equally unphysical adiabatic phase on large scales. The scheme with a fixed dissipation time of 10 Myr produces far too violent feedback at all resolutions, completely destroying the disc and giving rise to an unphysical pattern of star formation as gas is ejected from the system. This suggests our choice of parameters is incorrect, though we test a higher effective velocity dispersion threshold (100 km s$^{-1}$ instead of our fiducial 10 km s$^{-1}$) without a drastic change in results (see Appendix 2.C) and also note that our choice of parameters is not wildly different from others used in the literature at similar resolutions (see e.g. Teyssier et al. 2013; Rosdahl et al. 2017, though Dubois et al. 2015 determine lower values for the dissipation time). Our attempts to modulate the dissipation time with resolution ($t_{\text{diss}} = \Delta x/\sigma_{\text{FB}}$), suggested in Teyssier et al. (2013) as an alternative parametrization, do not converge with resolution, being similar to the fixed dissipation time at low resolution while essentially acting as a simple thermal dump at high resolution. No doubt both these schemes could be improved were we to spend more time tuning the parameters, but this would
not achieve our goal of finding a physically motivated model ideally free from adjustable parameters. We also note the concerns of Rosdahl et al. (2017) that their delayed cooling scheme trialled does not converge with their thermal dump when the adiabatic phase is resolved (as we find), suggesting that the scheme does not necessarily converge to the correct answer.

The most successful scheme explored is the mechanical feedback scheme. It suppresses star formation by similar factors across two orders of magnitude in mass resolution (though is slightly stronger at higher resolution), prevents the formation of highly dense clumps of gas, preserves the disc structure and agrees with observations of the Kennicutt-Schmidt relation (though the exact position on the relation has a resolution dependence caused by non-convergent outflow properties, discussed below). It also gives similar results to the classical schemes at the highest resolution, suggesting it is converging onto the correct answer. This latter feature was also noted in Rosdahl et al. (2017) and demonstrates that the ability of the scheme to converge on the final momentum input to the ISM per SN (as shown in Kimm & Cen 2014) translates into convergent behaviour for global properties. The mechanical feedback is slightly stronger than the classical feedback schemes because, even at this resolution, they are likely to still experience some overcooling at the highest density SN sites. The one area where the mechanical scheme does not converge is in outflows, which we discuss next.

2.4.2 Difficulties in outflow generation and the possible effects of missing physics

Having concluded in the previous section that the mechanical feedback scheme is the best amongst those explored, we choose to focus on it for this discussion. At the highest resolution, the scheme produces well developed multiphase outflows with appropriate mass loadings compared to observations and theory ($\beta_v \approx 1 - 10$), which is very encouraging (similar results are obtained by the classical schemes and the variable $t_{\text{diss}}$ delayed cooling at this resolution). What is not so encouraging, however, is that the outflows are considerably weaker at a mass resolution of 200 M$_{\odot}$ and practically non-existent with 2000 M$_{\odot}$, despite similar results with other galaxy properties. This also has the effect of moving the discs up the Kennicutt-Schmidt relation with decreasing resolution because the disc surface density is increased (though it should be noticed that these still match observations). Rosdahl et al. (2017) also report difficulties in driving outflows with mechanical feedback with a resolution similar to our lower resolutions simulations.
Such inefficient outflows could be caused by an oversimplified model of SN expansion. The mechanical feedback scheme treats the unresolved evolution of the SN remnant as expanding through a uniform medium. In reality, the ISM is likely to be porous due to a turbulent structure, containing low density channels through which gas accelerated by the SN can escape, leading to higher velocities. Haid et al. (2016) model this effect by considering the ISM surrounding the SN as a set of cones of different densities (randomly drawn from a log-normal distribution appropriate for the level of turbulence assumed) and use the results for a uniform medium within each cone. They find that momentum can be boosted by up to a factor of 2 in a low density environment. Our mechanical scheme already approximates this approach because it calculates the boost factor (eq. 2.21) for each neighbour cell independently, a point that is argued by Hopkins et al. (2018). However, we would caution that this assumes that the turbulent structure of the ISM is well resolved in the simulation, which is very unlikely to be the case, and will introduce a resolution dependence. Furthermore, the momentum boost measured in Haid et al. (2016) is weak; similar results are found in other studies with full 3D simulations (e.g. Iffrig & Hennebelle 2015; Martizzi et al. 2015; Kim & Ostriker 2015; Li 2015; Walch & Naab 2015), some of which find a slight negative impact on final momentum input versus the uniform case. However, while the final momentum input into the ISM may be only weakly affected by a turbulent medium, the amount of mass involved in the expansion and therefore the wind velocities reached can be altered. Kimm et al. (2015) trial a modification of their version of mechanical feedback in AMR simulations where they reduce the mass entrained from the host cell to 10% to replicate this effect, resulting in a greater suppression of star formation and higher mass loading factors. They state, however, that this fraction was somewhat arbitrarily chosen. Unless a physically motivated method of determining the fraction to be entrained based on unresolved structure was used, this could easily become just another tunable parameter. A final problem is that with a constant mass, as by definition imposed with a Lagrangian code, there is a minimum mass that can be momentum boosted. Even if we inject the correct momentum, the effects can be diluted if it is injected into too much mass resulting in lower velocity winds, effectively imposing a minimum resolution requirement.

Another potential cause of inefficient outflows experienced here could be the lack of other stellar feedback mechanisms. In particular, the ability of other feedback mechanisms to disrupt GMCs will enhance subsequent SN feedback. An obvious mechanism is that of photoionization (see e.g. Vázquez-Semadeni et al. 2010; Walch et al. 2012; Dale et al. 2014; Sales et al. 2014). Geen et al. (2015)
found that the final momentum input to the ISM by SN is increased when the surrounding medium has been preprocessed by photoionization feedback by forming an over-pressurised and lower density region in which the SN occurs. Kimm et al. (2017) modify their mechanical feedback prescription to include this momentum boost when they under-resolve the Strömgren sphere in their RHD simulations. Hopkins et al. (2012a, 2014) find that if they turn off radiative feedback (radiation pressure, photoionization and photoelectric heating), outflow mass loadings are reduced because GMCs are no longer efficiently disrupted prior to SN occurring (though the strength of the effect is dependent on the mass of the system). Simulating an isolated system similar to our fiducial model, Hu et al. (2017) find that while SNe are the dominant feedback mechanism, the inclusion of photoionization increases outflow rates by reducing the ambient density at SN sites. However, in their simulations, the inclusion of photoelectric heating reduces outflows because it reduces the SFR and therefore the number of SN occurring, while being unable to drive outflows itself.

As noted in Section 2.3.8, the choice of star formation prescription can also impact the effectiveness of feedback. We found that increasing the star formation efficiency parameter by a factor of 10 led to stronger outflows (and the destruction of the disc) because the SFR was initially higher and gas could not reach high densities before SN occurred, leading to a sudden, strong burst of efficient feedback. Instead, increasing the threshold density had only a marginal impact on the SFR, but produced weaker outflows because SN occurred in slightly denser environments. The adoption of other feedback mechanisms would probably mitigate this effect. It is also worth noting that we have only tested changes of parameters to our simple star formation prescription. More complex prescriptions may rely on the selection criteria of star forming gas rather than a efficiency parameter (see e.g. Hopkins et al. 2013, 2014, and subsequent papers, which use an efficiency of 100%, but require gas to be self-gravitating, self-shielding and very dense). Alternatively, it has been suggested that while the globally averaged star formation efficiency may be on the order of a few percent, small scale efficiencies vary based on the local properties of the ISM (see e.g. Krumholz & McKee 2005; Padoan & Nordlund 2011; Hennebelle & Chabrier 2011; Federrath & Klessen 2012). In line with this, a star formation prescription could adopt a variable efficiency (see e.g. Kimm et al. 2017). Such schemes are likely to impact the distribution of gas densities by allowing high density non-star forming gas to exist, as well as impacting SN feedback effectiveness by altering the clustering properties of stars in both space and time.
2.5 Conclusion

Using an isolated disc galaxy setup and a new implementation of star formation and SN feedback in the moving mesh code AREPO, we tested several SN feedback prescriptions commonly found in the literature and assessed their impacts on a variety of galaxy metrics, paying particular attention to how well they converge as a function of resolution. The bulk of our simulations were of a $10^{10} \, M_\odot$ system, although simulations were carried out of systems an order of magnitude lower and higher in mass. In order to test the convergent properties of the feedback schemes with resolution, simulations were carried out with resolutions of $2000 \, M_\odot$, $200 \, M_\odot$ and $20 \, M_\odot$. The schemes tested were designed to be used in isolated galaxy or cosmological zoom-in simulations, using individually time resolved SN events. Specifically, we investigated ‘classical’ dumps of thermal and/or kinetic energy, two parametrizations of delayed cooling and a mechanical feedback scheme which injects the correct amount of momentum relative to the stage of the SN remnant evolution resolved.

Without feedback, our simulations produce a highly clumpy disc and overproduce the mass of newly formed stars. As expected, the ‘classical’ feedback schemes overcool at all but the highest resolution. The delayed cooling schemes tested are far too strong, unphysically destroying the disc. We note that we could tune these simulations more carefully to avoid this effect, but because we wish to avoid adjustable parameters as much as possible we do not consider these schemes to be suitable for our purpose. Our mechanical scheme is the best tested, suppressing star formation by similar factors at all three resolutions, preventing the formation of highly dense clumps of gas, agreeing with observations of the Kennicutt-Schmidt relation while also preserving the disc structure. It also produces similar results to the ‘classical’ schemes at $20 \, M_\odot$ resolution, suggesting it is converging onto the physically correct results.

At the highest resolution our mechanical scheme produces multiphase outflows with reasonable mass loading factors relative to observations and theory ($\beta_v \approx 1 - 10$), as do the ‘classical’ schemes at the highest resolution. However, we struggle to produce outflows at lower resolution. This may be due to an oversimplification of the way in which we model SN remnant evolution, for example failing to adequately account for the unresolved porous structure of the ISM. The situation may also be improved by the inclusion of other forms of stellar feedback that are able to preprocess the ISM and enhance the ability of the SN feedback to drive galactic winds. In addition, alternative star formation prescriptions that aim to better
capture small scale star formation physics will impact the effectiveness of feedback by altering the clustering properties of SNe (both in space and time). It should also be noted that there exists some minimum resolution requirement for the driving of outflows with individually time resolved SN because the injection of momentum (even if it is the physically correct amount) into too much mass will result in unphysically slow gas velocities.

Finally, it is worth pointing out that the resolution requirements for the mechanical feedback scheme to work well in terms of outflows is within the reach of next generation cosmological zoom-in simulations (at least for low mass systems). This will allow us to explore realistic SN feedback in a full cosmological environment, self-consistently taking into account the circulation of complex gas flows all the way from the cosmic web to the ISM. The adoption of a more accurate star formation prescription in concert with the inclusion of other forms of stellar feedback in such simulations may ultimately help us unveil what shapes star formation in low mass systems.

2.A Non-thermal pressure floor

As described in Section 2.2.2 we impose a non-thermal pressure floor to avoid artificial fragmentation that may occur when the Jeans length is not properly resolved. Setting a minimum pressure using equation (2.1) ensures that the resulting Jeans length is always resolved by at least $N_J$ cells. Truelove et al. (1997) suggests that, at a minimum, the Jeans length must be resolved by at least 4 cells. This criteria is widely adopted in gravitational hydrodynamic simulations, but a variety of values can be found in the literature. The choice of $N_J$ is non-trivial: too low and artificial fragmentation will occur, too high and the formation of small (physical) structures that would otherwise be resolved is suppressed. Fig. 2.16 shows the effect of various choices of $N_J$ on the morphology of our fiducial galaxy model at 100 Myr with no feedback with $1000 M_\odot$ resolution. It can be seen that with no pressure floor ($N_J = 0$) the disc fragments into multiple, small, high density clumps. At the other extreme, enforcing resolution of the Jeans length by 16 cells ($N_J = 16$) results in the washing out of all small structure save some weakly defined spiral arms. Other choices of $N_J$ in between these values results in a corresponding sliding scale of morphologies. We can reasonably confidently assume that the adoption of $N_J = 16$ results in a morphology that is over-smoothed. Unfortunately, it is not so easy to say what the correct lower limit of $N_J$ is. Determining exactly when the onset of artificial fragmentation occurs is a non-trivial problem that is beyond the scope of
Fig. 2.16 Face-on gas density projections of simulations after 100 Myr with varying values of $N_J$ used to determine the artificial pressure floor, where $N_J$ is the number of cells by which the Jeans length must be resolved. Each simulation is carried out with no feedback and at 1000 $M_\odot$ resolution using our fiducial galaxy model. While imposing no floor results in artificial fragmentation on the smallest scales (top left panel), the $N_J = 16$ run (bottom right) washes out the physical structures present in the disc.

We attempt to crudely quantify the degree of fragmentation in Fig. 2.17 by plotting the mass of gas above some density and the mass of newly formed stars for our various choices of $N_J$. Mass of gas above $100 \text{ cm}^{-3}$ decreases by well over an order of magnitude across the range of $N_J$ probed, with even more dramatic increases when examining higher densities and, perhaps more worryingly, mass in new stars. It can be seen that a definite transition occurs from a regime of suppression of high density gas to a regime where gas may reach these densities, though this transition happens at different values of $N_J$ depending on the density examined. For example, for densities above $100 \text{ cm}^{-3}$ the transition occurs at $N_J \sim 10$, while for densities above $10^3 \text{ cm}^{-3}$ and $10^4 \text{ cm}^{-3}$ the transitions occurs at $N_J \sim 8$ and $\sim 6$, respectively. This transition could represent the transition between an artificially fragmenting regime to a stabilised regime, but without a more careful determination of what ‘artificial fragmentation’ is, one could just as easily state that it merely marks the transition from over-smoothing of structure to
Supernova feedback in isolated simulations of galaxy formation

Fig. 2.17 The mass above a given density (or stellar mass, grey dashed curve) as function of $N_J$ at 100 Myr. Each value is measured from a simulation with no feedback at 1000 M$_\odot$ resolution using our fiducial galaxy model. Note that there is a clear suppression of gas mass above a given density for densities above 100 cm$^{-3}$ which motivates our choice of the reasonable $N_J$ value to adopt, but for the stellar mass no such trend exists.

Typically, from empirical results (e.g. Truelove et al. 1997), one assumes that there is a single ‘correct’ value of $N_J$ that should be used to avoid artificial fragmen-
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Fig. 2.18 Face-on gas density projections after 100 Myr comparing our fiducial choice of \( N_J = 8 \) for a 1000 M\(_\odot\) resolution simulation with the same value for a 100 M\(_\odot\) resolution simulation (i.e. assuming that the ‘correct’ choice for \( N_J \) is resolution independent) and \( N_J = 17.2 \) (i.e. such that the Jeans length is resolved by the same physical scale between resolutions). Each simulation is with no feedback using our fiducial galaxy model.

The feedback and that the value is resolution independent. In other words, there exists some minimum required number of cells to correctly resolve fragmentation. Using a fixed value of \( N_J \) then allows fragmentation to occur on smaller scales as the minimum resolvable length decreases as resolution is increased. While this is often a desirable behaviour, particularly when concerned with ISM properties on the edge of the resolution limit, it necessarily leads to divergent galaxy properties. Instead, scaling \( N_J \) such that the minimum resolved Jeans length corresponds to the same physical scale at all resolutions results in convergent morphologies. In Fig. 2.18, we compare the 1000 M\(_\odot\) resolution simulation with \( N_J = 8 \) (as in Fig. 2.16) with two simulations with a resolution of 100 M\(_\odot\). One uses a value of \( N_J = 8 \), the other uses a value of \( N_J = 17.2 \) (i.e. scaled with mass resolution such that the minimum Jeans length is the same at both resolutions). It is clear that using the same value of \( N_J \) results in very different morphologies, whereas the adoption of a higher value with better resolution results in a similar morphology. This is not to say that the 1000 M\(_\odot\) resolution, \( N_J = 8 \) and 100 M\(_\odot\) resolution, \( N_J = 17.2 \) results are more ‘correct’ than the other, since as previously mentioned, it is difficult to determine when artificial fragmentation occurs. However, for the purposes of this paper where we are primarily concerned about the effects of differing SN feedback implementations, particularly examining the role of the resolution adopted on the feedback, we find that it is advantageous to enforce approximately similar galaxy
evolutions in the no feedback case across our resolutions by scaling the value of $N_J$ with resolution as described in Section 2.2.2. This we broadly achieve, with the no feedback simulations at different resolutions producing same amount of stars within a factor of a few and having comparable morphologies (see Figs. 2.2, 2.21 and 2.22). It is worth highlighting that we are not unique in our choice to scale $N_J$ with resolution (see for example Rosdahl et al. 2015, 2017).

When we run simulations without a pressure floor but with feedback, the effect is similar to increasing the star formation efficiency parameter (see Fig. 2.14), with a sudden burst of high SFR followed by extremely strong feedback that largely destroys the disc and quenches star formation. It is possible that if we included other stellar feedback mechanisms (stellar winds, radiation pressure, photoionization) that are active in the intervening time between star formation and SN occurrence it might be possible to keep gas from entering the regime where it is vulnerable to artificial fragmentation, thus removing the need for a pressure floor. Alternatively, a more complex star formation criteria that identifies fragmenting gas could also circumvent the issue (e.g. Hopkins et al. (2017) argue that the Jeans unstable gas should be turned into stars rather than using a pressure floor), but if the Jeans length is significantly under-resolved this could result in the spurious boosting of SFRs. We conclude this section by remarking that, on the whole, when artificial pressure floors are adopted, the motivation behind the choice of parameters is often not clear. Given the strong dependence of results on this choice, we suggest that this is an issue that needs to be addressed in more detail in future work.

2.B SPH-like kernel weighting versus explicitly isotropic weighting scheme for SN feedback

As mentioned in Section 2.2.4, we have found that under certain conditions, the use of a simple SPH-like kernel-based weighting scheme for distributing feedback quantities (mass, metals, energy and momentum) into the gas local to the SN can result in significant violations of the desired isotropic distribution. Because such a weighting scheme preferentially injects into denser regions where there are more cells, if there is a strong density gradient in a particular direction injection of feedback quantities will be injected perpendicular to the gradient. For example, in the case of a SN occurring in a thin disc, more resolution elements lie in the disc plane than lie above and below it so feedback quantities will be preferentially injected into the disc plane. The situation is exacerbated by poor resolution and by the use of efficient momentum-based feedback schemes (thermal injection based
Fig. 2.19 Face-on gas density projections of simulations at 100 Myr comparing the use of an SPH-like mass weighting scheme for distributing SNe mass, energy and momentum with our explicitly isotropic weighting scheme described in Section 2.2.4. The simulations have a resolution of 2000 $M_\odot$, use our fiducial galaxy model with a modified form of our mechanical feedback (see main text). SPH-like weighting leads to unphysical shells propagating through the disc, sweeping up most of its mass. Our isotropic weighting scheme avoids this numerical issue correctly coupling the SN ejecta to the surrounding gas regardless of its density.

schemes can mitigate the situation slightly since heated cells will tend to expand along the path of least resistance). This can result in the unphysical driving of expanding shells through the disc plane, with little ejecta going in the vertical direction (see also Hopkins et al. 2017; Hopkins et al. 2018).

Fig. 2.19 demonstrates this effect. We compare a simulation using a standard SPH-kernel based mass weighting scheme for distributing feedback quantities with our explicitly isotropic weighting scheme as described in Section 2.2.4. For numerical reasons, we cannot use our full mechanical feedback scheme with an SPH-like scheme, so we use a hybrid of our kinetic and mechanical feedback schemes for this
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comparison; we inject $2.41 \times 10^5 \text{ km s}^{-1} \text{ M}_\odot$ of momentum per SN, corresponding to the final momentum of a SN occurring in gas of density $100 \text{ cm}^{-3}$ and metallicity $0.1 \ Z_\odot$ as calculated using equation (2.22). The simulations are of our fiducial galaxy at $2000 \text{ M}_\odot$ resolution and the projections shown in Fig. 2.19 are at 100 Myr. The SPH-like weighting scheme sweeps the disc mass into a thin expanding ring, whereas our isotropic scheme prevents this occurring. This scenario is where the effect is most noticeable, but it is still present to some extent at higher resolutions.

It should be noticed that switching to a volume weighting scheme rather than mass weighting does not have much of an effect. For a reasonable neighbour number (32 - 64, as used with a cubic spline kernel) most identified neighbours will be in the disc plane. If cells are found within the smoothing length containing the neighbours that lie above the plane of the disc, the extra weighting they will receive for being of a larger volume (because they are less dense) is likely to be subdominant compared to the ‘penalty’ they receive for being furthest away from the star particle.

2.C Other delayed cooling parameters

As our fiducial parameters for the delayed cooling with fixed dissipation time we have adopted $t_{\text{diss}} = 10 \text{ Myr}$ and $\sigma_{\text{FB}} = 10 \text{ km s}^{-1}$, as used in Teyssier et al. (2013). As noted above, with our galaxy models at all resolutions explored, this feedback scheme appears to be very strong relative to our other schemes and produces unphysical results. We therefore tried a higher threshold velocity dispersion, $\sigma_{\text{FB,threshold}} = 100 \text{ km s}^{-1}$, as used in Rosdahl et al. (2017). Fig. 2.20 shows the effect of using these parameters on the star formation rate and stellar masses with our fiducial galaxy model at all three of our resolutions. The SFRs are not suppressed to the same degree with this weaker feedback, with final new stellar mass being approximately a factor of 2 larger at all resolutions. However, the feedback still destroys the disc in the same manner as our fiducial simulations (though gas returns to the centre, resulting in a second burst of star formation). As mentioned above, with a more careful approach to tuning these parameters, we could perhaps arrive at a less aggressive scheme.

2.D Other resolutions

This appendix contains results for our lower resolution simulations for comparison to the figures in the main text that show our highest resolution simulations. Figs. 2.21 and 2.22 show face-on and edge-on density projections of gas and newly
Fig. 2.20 Newly formed stellar mass (top) and SFRs (bottom) for simulations at all three resolutions for runs with delayed cooling using our fiducial threshold of 10 km s$^{-1}$ and a higher value of 100 km s$^{-1}$. The results are not very sensitive to the change of $\sigma_{FB,\text{threshold}}$ and while the SFRs are suppressed somewhat less with the choice of the higher value, the disc is still largely disrupted after the first peak in SFR.

formed stars after 250 Myr (see Fig. 2.2 for the highest resolution case). Fig. 2.23 shows density, temperature and metallicity slices at 250 Myr to demonstrate how outflow properties change with resolution (see Fig. 2.8 for the highest resolution case).
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**Fig. 2.21** Projections of gas and new stars formed for our different feedback simulations at 250 Myr for 2000 M$_\odot$ resolution runs. The mixed feedback simulation is not shown as the results are similar to the thermal and kinetic feedback simulations. The morphology of the no feedback simulation disc is similar to the highest resolution case (see Fig. 2.2), with a highly clumpy distribution of gas and newly formed stars, although the structure is less well defined at this resolution (particularly in the stellar component). At this resolution, the classical feedback schemes overcool, so they produce similar morphologies to the no feedback simulation. Both delayed feedback schemes are too powerful, disrupting the gas disc and producing ring-like structures of newly formed stars. The mechanical feedback scheme is able to suppress the formation of dense clumps without destroying the disc. The resulting morphology is similar to the high resolution simulation, although it is not as well defined.
Fig. 2.22 Projections of gas and new stars formed for our different feedback simulation at 250 Myr for 200 M⊙ resolution runs. The mixed feedback simulation is not shown as the results are similar to the thermal and kinetic feedback simulations. The morphology of the no feedback simulation disc is similar to the highest and lowest resolution cases (see Figs. 2.2 and 2.21). As in the lowest resolution simulation, the classical feedback schemes overcool and produce a similar clumped morphology to the no feedback simulation. The delayed feedback schemes remain too powerful, disrupting the gas disc, although the scheme with variable \( t_{\text{diss}} \) is weaker. The mechanical feedback scheme produces a similar morphology to the lower and higher resolutions simulations.
Unsurprisingly, the no feedback simulations do not produce outflows. Only the delayed cooling schemes drive outflows at both resolutions, ejecting the majority of material from the centre of the system. At the 200 M\(_\odot\) resolution, the classical feedback schemes produce outflows that reach a short distance above the disc. They are highly metal enriched because of the large number of SNe driving the outflows, due to the inefficiency of the feedback caused by overcooling. The mechanical feedback is able to drive a modest outflow at the 200 M\(_\odot\) resolution, though the outflow peaked \(\sim 100\) Myr previously resulting in material returning to the disc in a galactic fountain (see Figs. 2.6 and 2.7). Figure continued overleaf.
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Fig. 2.22 (Cont.)
3. Cosmological simulations of dwarfs

The need for ISM physics beyond SN feedback alone

The dominant feedback mechanism in low mass haloes is usually assumed to take the form of massive stars exploding as supernovae (SNe). We perform very high resolution cosmological zoom-in simulations of five dwarf galaxies to $z = 4$ with our mechanical SN feedback model. This delivers the correct amount of momentum corresponding to the stage of the SN remnant evolution resolved, and has been shown to lead to realistic dwarf properties in isolated simulations. We find that in 4 out of our 5 simulated cosmological dwarfs, SN feedback has insufficient impact resulting in excessive stellar masses, extremely compact sizes and central supersolar stellar metallicities. The failure of the SN feedback in our dwarfs is physical in nature and is the result of the build up of very dense gas in the early universe due to mergers and cosmic inflows prior to the first SN occurring. We demonstrate that our results are insensitive to resolution (provided that it is high enough), details of the (spatially uniform) UV background and reasonable alterations to the star formation prescription. We therefore conclude that other physical processes, such as additional forms of stellar feedback, are required to prevent the excessive build up of dense gas hence allowing SNe to regulate dwarf properties.

3.1 Introduction

While they may be the least massive and luminous systems in our universe, understanding the origin of dwarf galaxy properties represents a key step in developing and testing theories of galaxy formation and cosmology. From the perspective of cosmology, abundances and structural properties of low mass haloes present an important observational test of the $\Lambda$CDM model, but the effects of baryonic physics (to which dwarfs are highly susceptible because of their small potential wells) can make it difficult to make robust predictions. Meanwhile, from the point of view of galaxy formation, these ‘messy’ baryonic processes are interesting in their own right, as well as providing insight into the reionization history of the universe and the enrichment of the intergalactic medium (IGM) with metals.

Examining the issue of dwarf abundances, there is a substantial offset between the predicted dark matter halo abundance from numerical simulations and the
observed galaxy stellar mass function (for recent work see e.g. Behroozi et al. 2013; Moster et al. 2018), indicating that dwarfs must be more than an order of magnitude less efficient at forming stars than Milky Way-sized haloes. This is also posited as a solution to the so called ‘missing satellites problem’, where the observed number of Local Group satellites is at odds with the substantially larger number of dark matter haloes predicted by cosmological N-body simulations (see e.g. Moore et al. 1999; Klypin et al. 1999; Diemand et al. 2008; Springel et al. 2008; Koposov et al. 2009; Rashkov et al. 2012; Sawala et al. 2016).

It has been suggested for some time that low mass haloes should have their star formation efficiency suppressed by two primary processes: SN feedback (e.g. Larson 1974; Dekel & Silk 1986; Mori et al. 2002; Governato et al. 2007) and cosmic reionization (e.g. Efstathiou 1992; Bullock et al. 2000; Dijkstra et al. 2004; Kravtsov et al. 2004b; Madau et al. 2008). Evidence that these structures do in fact exist but are relatively dark has been bolstered recently by detections of local ultra-faint dwarfs (e.g. Koposov et al. 2015; Laevens et al. 2015; Martin et al. 2015; Kim et al. 2015).

As well as influencing abundances of low mass haloes, baryonic physics has also been invoked to solve structural discrepancies between dark matter simulations and observations. One such discrepancy is often termed the ‘cusp-core controversy’. Within the ΛCDM paradigm, dark matter-only simulations systematically predict steep inner density profiles for these low mass haloes, but some observations suggest that they may instead contain low density cores (see e.g. Moore 1994; Flores & Primack 1994; de Blok & Bosma 2002; Walker & Peñarrubia 2011). While SN feedback has been widely invoked in hydrodynamical simulations in an attempt to generate cored density profiles, there is still no consensus in the literature as some groups find only cuspy profiles (e.g. Vogelsberger et al. 2014; Sawala et al. 2016), while others find various levels of cored profiles with different trends as a function of halo mass or redshift (e.g. Navarro et al. 1996; Gnedin & Zhao 2002; Read & Gilmore 2005; Mashchenko et al. 2008; Governato et al. 2010; Pontzen & Governato 2012; Di Cintio et al. 2014; Oñorbe et al. 2015; Fitts et al. 2017). The level of success in transforming dark matter cusps into cores seems closely related to the degree of burstiness of SN feedback, which could also affect the mass-loading of galactic outflows and the early enrichment of the IGM.

It is perhaps at some level unsurprising that the properties of simulated dwarfs predicted by different groups are at variance as very different sub-grid models for star formation, SN feedback and wind launching are adopted, in addition to results often being rather sensitive to the numerical resolution of the simulations. This
is however clearly unsatisfactory if we are to understand at a more fundamental level how SN feedback operates in dwarf galaxies, and even more so if we are to derive robust constraints on the nature of dark matter, using observed dwarfs as near-field cosmology probes.

Recently, there have been several theoretical works (e.g. Hopkins et al. 2014; Kimm & Cen 2014; Kim & Ostriker 2015; Walch & Naab 2015; Martizzi et al. 2015, but see also earlier work by Cioffi et al. 1988; Thornton et al. 1998) aiming at quantifying the correct momentum injection at a given SN remnant stage as a function of local ISM properties (such as the gas density, metallicity and porosity), based on analytical calculations or small scale simulations of individual SN explosions. These studies are particularly useful because, in principle, they allow us to impart the appropriate momentum per SN to the ISM (without the use of tunable parameters), even when the Sedov-Taylor phase of the SN remnant evolution is not properly resolved, which is often the case in galaxy formation simulations. We have explored this type of SN injection, often dubbed ‘mechanical feedback’, in an extensive series of simulations of isolated dwarf galaxies (Chapter 2), finding that it results in realistic and well converged star formation rates and dwarfs morphologies over two orders of magnitude in mass resolution explored. However, to naturally produce mass-loaded, multiphase outflows, the gas mass resolution needed to be on the order of a few tens of $M_\odot$ at least.

Thankfully, even though these resolution requirements are quite daunting, they are achievable in full cosmological simulations provided that one dwarf is simulated at a time with a zoom-in technique. Hence, the aim of this work is to explore the mechanical SN feedback scheme in fully self-consistently formed dwarfs without tuning any parameters in order to understand if it leads to realistic dwarf properties once the cosmological gas inflows and mergers are taken into account. We explore this by randomly selecting five dwarfs with virial masses between $\sim 2-6 \times 10^9$ $M_\odot$ at $z = 4$ which reside in different environments and have different assembly histories.

### 3.2 Methodology

Our numerical scheme is essentially the same as that described in Chapter 2, but we summarise the salient details here. We carry out our simulations with the moving-mesh codeAREPO (Springel 2010) which solves hydrodynamics on an unstructured Voronoi mesh. Gravity is included using a hybrid TreePM scheme.

In this work, we include radiative cooling as in Vogelsberger et al. (2013). Primordial heating and cooling rates are calculated using cooling, recombination
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and collisional rates provided by Cen (1992) and Katz et al. (1996). Metal-line cooling to 10 K is obtained from lookup tables containing rates precalculated from the photoionization code CLOUDY. We include a redshift dependent, but spatially homogeneous UV background from Faucher-Giguère et al. (2009), although it is only turned on from $z = 9$ to approximate the latest Planck measurement of optical depth to reionization (Planck Collaboration XIII 2016). We adopt the density based self-shielding prescription of Rahmati et al. (2013) to attenuate the UV background in dense gas.

We include a non-thermal pressure floor to prevent artificial fragmentation in the event of under-resolving the Jeans length (see e.g. Truelove et al. 1997). To ensure that the Jeans length is resolved by $N_J$ cells, this takes the form

$$P_{\text{min}} = \frac{N_J^2 \Delta x^2 G \rho^2}{\pi \gamma},$$  \hspace{1cm} (3.1)

where $\Delta x$ is the cell diameter, $\rho$ is the gas density and $\gamma = 5/3$ is the adiabatic index. We adopt $N_J = 8$ in this work. A detailed discussion of the effects of adopting a pressure floor can be found in Chapter 2 (see also Section 3.4.2).

Gas above a density threshold of $n_{\text{SF}}$ is assigned a star formation rate density according to a simple Schmidt law,

$$\dot{\rho}_* = \epsilon_{\text{SF}} \frac{\rho}{t_{\text{ff}}},$$  \hspace{1cm} (3.2)

where $\rho$ is the gas density, $\epsilon_{\text{SF}}$ is some efficiency and $t_{\text{ff}} = \sqrt{3\pi/32G\rho}$ is the free-fall time. We use a fiducial value of $n_{\text{SF}} = 10 \text{ cm}^{-3}$ and $\epsilon_{\text{SF}} = 1.5\%$ (chosen to match observed efficiencies in dense gas, see e.g. Krumholz & Tan 2007, and references therein). We also examine the effect of varying these parameters in Section 3.4.2. Using these rates, gas cells are then stochastically converted into star particles (collisionless particles representing single stellar populations). Star particles inherit the metallicity of the gas from which they were formed.

For each star particle, we obtain a SN rate, $\dot{N}_{\text{SN}}$, as a function of age and metallicity precalculated using STARBURST99 (Leitherer et al. 1999) assuming a Kroupa (2002) IMF. The number of SNe that occur in a timestep is then drawn from a Poisson distribution with a mean of $\bar{N}_{\text{SN}} = \dot{N}_{\text{SN}} \Delta t$, where $\Delta t$ is the timestep. In order to individually time resolve SNe, we impose a timestep limiter for star particles to ensure that $\bar{N}_{\text{SN}} \ll 1$.

\footnote{However, we have performed extra simulations where the UV background follows Faucher-Giguère et al. (2009) exactly (switching on at $z = 11.7$) and find that it does not change our results in any appreciable way.}
When a SN occurs, mass, metals, energy and momentum are coupled to the gas cell containing the star particle (the host cell) and its neighbours (all those that share a face with the host cell). Feedback quantities are distributed to the gas cells using an explicitly isotropic weighting scheme in the rest frame of the star particle (details in Chapter 2, see also Hopkins et al. 2018) in order to avoid spurious numerical effects that can arise when using a simple kernel (mass) weighted approach to nearest neighbours due to the increased relative number of resolution elements present in dense gas. The ejecta mass, \( m_{\text{ej}} \), deposited per SN is \( 10 M_\odot \), of which \( 2 M_\odot \) is in metals, with an energy of \( 10^{51} \) ergs. In simulations designated ‘no feedback’, mass and metals are returned, but no feedback energy/momentum is deposited. In runs with full SN feedback, we adopt the mechanical feedback scheme described in Chapter 2 (see also Hopkins et al. 2014, 2017; Hopkins et al. 2018; Kimm & Cen 2014; Kimm et al. 2015; Martizzi et al. 2015). This aims to deposit the correct amount of momentum corresponding to the stage of the SN remnant evolution resolved (dependent on the local gas density and metallicity).

When analysing simulations, we use the halo finder SUBFIND (Springel et al. 2001; Dolag et al. 2009) to calculate halo properties. We adopt the convention of considering friends-of-friends (FOF) groups as the primary dark matter halo and subhaloes as galaxies within the halo (unless otherwise stated, we only consider centrals). For halo virial masses, we use the definition of Bryan & Norman (1998) and for galaxy stellar masses we use the mass contained within twice the radius that contains half the total subhalo stellar mass associated with the group. We use the code SUBLINK (Rodriguez-Gomez et al. 2015) to construct merger trees and track haloes/subhaloes throughout the simulations. We follow the branch of the merger tree with the most mass behind it for our analysis except where otherwise mentioned. We adopt a Planck Collaboration XIII (2016) cosmology throughout this work. Unless otherwise stated, all units are in proper coordinates.

3.3 Simulations

3.3.1 Initial conditions and simulation details

The process for creating cosmological ‘zoom-in’ initial conditions is as follows. First, a coarse resolution dark matter only simulation of a large, periodic cosmological volume is run to a target redshift, \( z_{\text{target}} \). Dark matter haloes of interest are identified in the \( z_{\text{target}} \) output of this simulation and are resimulated at a higher resolution with a ‘zoom-in’ technique. Gas is then added to the initial conditions.
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Fig. 3.1 Density projections of the large scale environment around our target haloes in the coarse dark matter only simulation at \( z = 4 \). The target haloes are marked with green ticks for ease of identification and the virial radius is marked with a green circle. We deliberately select the haloes from a variety of environments ranging from void-like regions to rich filaments.

by splitting the particles into dark matter and gas mesh generating points according to the cosmic baryon fraction (although we also carry out dark matter only zoom-in simulations).

We use the code MUSIC (Hahn & Abel 2011) to generate the initial conditions (at \( z = 127 \)) for both the coarse and zoom-in simulations. Dwarfs 1 and 2 are selected at \( z = 0 \) from 10 cMpc h\(^{-1}\) coarse boxes with a resolution of 256\(^3\) particles (giving a particle mass of \( 7.47 \times 10^6 \) M\(_\odot\)). In the coarse simulation, their virial masses at \( z = 0 \) are \( 1.04 \times 10^{10} \) M\(_\odot\) and \( 1.12 \times 10^{10} \) M\(_\odot\) with virial radii of 62.0 kpc and 64.5 kpc, respectively. The selection regions at \( z = 0 \) are a sphere of radius 736 kpc for dwarf 1 and a sphere of radius 295 kpc for dwarf 2. However, for the purposes of this work, we carry out our analysis up until \( z = 4 \), at which point their masses are \( 2.82 \times 10^9 \) M\(_\odot\) and \( 3.11 \times 10^9 \) M\(_\odot\) (note that in the zoom-in simulations, the final mass varies due to the effects of baryonic physics and the higher resolution).

Dwarfs 3, 4 and 5 are selected at \( z = 4 \) from a 20 cMpc h\(^{-1}\) box with a resolution of 512\(^3\) (i.e. the same mass resolution as the boxes used for dwarfs 1 and 2). The masses of dwarfs 3 and 4 at \( z = 4 \) are \( 2.56 \times 10^9 \) M\(_\odot\) and \( 2.51 \times 10^9 \) M\(_\odot\) with virial radii of 8.86 kpc and 8.78 kpc. The selection regions at \( z = 4 \) are spheres of radius 44.1 kpc. In the coarse simulation, we identify a fifth halo with a virial mass of \( 1.00 \times 10^{10} \) M\(_\odot\) and a virial radius of 13.96 kpc, which we resimulate with a zoom-in region of 88.3 kpc. In the subsequent zoom-in simulation, this region actually contains two separate haloes of \( \sim 6 \times 10^9 \) M\(_\odot\), sufficiently separated as to be considered independent. We take the larger of these two haloes to be the focus of our analysis, referring to it as dwarf 5. Fig. 3.1 shows dark matter
density projections of the large scale region around the target haloes in the coarse simulations at $z = 4$. Dwarfs 1, 2 and 3 are in relatively low density environments, 4 is in a more crowded filament region, while 5 is a larger system in a very crowded filament.

Our fiducial simulations increase the number of resolution elements in the zoom-in region by a factor of $16^3$ giving dark matter particle and target gas cell masses of $1536\, \text{M}_\odot$ and $287\, \text{M}_\odot$, respectively (we also run a simulation of dwarf 1 with a higher resolution of $191\, \text{M}_\odot$ in dark matter and $35.9\, \text{M}_\odot$ in gas for the purposes of testing convergence). The refinement/derefinement scheme in AREPO keeps gas cell masses within a factor of 2 of the target mass. Because star particles are formed by converting gas cells, this also corresponds to the initial star particle mass (prior to mass loss from feedback). We use comoving gravitational softenings of $0.129\, \text{ckpc}$ for the high resolution dark matter particles, gas cells\(^2\) and star particles. For dwarfs 3, 4 and 5 the softening is held at its $z = 6$ proper length of $18.4\, \text{pc}$ from that redshift onwards, although this makes very little practical difference. Because our simulations do not include the necessary physics (such as molecular cooling) to resolve Population III stars and the first enrichment of the ISM, we initially assign gas in the simulation a metallicity of $10^{-4}\, \text{Z}_\odot$. For each dwarf, we carry out a simulation to $z = 4$ with no feedback and with SNe (due to computational expense, we run the no feedback dwarf 5 simulation to $z = 5.5$). Section 3.4 contains details of additional simulations of dwarf 1 carried out with various modifications to our fiducial parameters to test convergence.

3.3.2 Results

Fig. 3.2 shows density projections of the dark matter, gas and stars for our simulated dwarfs at $z = 4$ (dwarf 5 is shown at $z = 5.5$ to allow comparison between the runs with and without feedback). A variety of morphologies are present. Runs without feedback tend to produce highly compact gas and stellar discs. Recent mergers can give rise to warped structures, for example in dwarfs 1 and 4. With feedback, dwarfs 1, 2 and 5 also feature compact stellar discs similar to the no feedback simulations, although their orientation has changed. The gas morphology is more obviously changed, with more irregular and diffuse structure. In dwarfs 3 and 4, feedback has made a significant impact on the stellar structure, with significantly lower surface densities and the absence of a well defined disc. In dwarf 3, most of the gas has been cleared away, leaving a small dense core, while the

\(^2\)For gas cells, the softening is calculated as the maximum of either this fixed softening value or 2.5 times the cell radius.
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Fig. 3.2 Density projections, from left to right: dark matter (shown here for the no feedback simulations, although the equivalent plots for runs with SNe are similar), gas for simulations without SNe, gas with SNe, stars without SNe and stars with SNe. Each row corresponds to a different dwarf. The gas and stellar projections are centred on the central galaxy of the halo. Dwarfs 1, 2, 3 and 4 are shown at $z = 4$, while dwarf 5 is shown at $z = 5.5$ to allow comparison to the curtailed no feedback simulation. While SN feedback significantly alters morphologies, particularly in the case of dwarf 4, in most cases centrally condensed baryon concentration persists.

feedback has almost completely evacuated the gas from dwarf 4.

Fig. 3.3 shows the stellar mass to halo mass ratio of the central galaxies formed in the simulations, normalised by the cosmic baryon fraction, as a function of
Fig. 3.3 Stellar mass to halo mass ratio as a function of halo mass for our various simulations at $z = 10, 8, 6$ and 4. Halo mass here is defined as in Bryan & Norman (1998). We express the stellar mass to halo mass ratio (which can be considered the integrated baryon conversion efficiency) as the mass of stars formed in the central galaxy (defined as the most massive subhalo) divided by the product of the halo mass and the cosmic baryon fraction. Open symbols indicate simulations without SNe. In the case of dwarf 4, we plot both the progenitor haloes of the final halo prior to their major merger at $z = 5.5$. At $z = 4$, while the two haloes have merged, the two central galaxies of the progenitors have not yet merged into a single subhalo (see main text); nonetheless, we use the sum of the stellar mass in both of these galaxies to compute the star formation efficiency. We also indicate results from abundance matching as in Behroozi et al. (2013) and Moster et al. (2018) with shaded regions, although it should be noted that at such low halo masses the relations are heavily extrapolated. Note that with exception of dwarf 4, all of our simulated dwarfs are in large disagreement with the abundance matching extrapolations.

halo mass, for 4 redshifts ($z = 10, 8, 6, 4$). We also plot empirically derived abundance matching results from Behroozi et al. (2013) and Moster et al. (2018) for comparison, although we have heavily extrapolated the results to reach this
mass range so they should be treated with caution. However, even with this caveat in mind, it can be seen that the majority of our simulated galaxies massively overproduce stars, lying several orders of magnitude above the abundance matching relations at all four redshifts. This is true for all simulations without SNe, where typically $10 - 60\%$ of the available baryons (taking that to be $f_b M_{\text{halo}}$) has been converted into stars, with variation of only a factor of a few between $z = 10 - 4$. With feedback, there are mixed results. Dwarf 1 produces almost identical stellar to halo mass ratios with and without feedback at all redshifts, with only marginal suppression of star formation by $z = 4$. Similarly, in dwarf 5 feedback has little impact on the evolution of stellar mass. For dwarf 2, at $z = 10$, the ratio is about an order of magnitude lower in the run with feedback than without (although still somewhat high). However, the ratio increases with decreasing redshift and by $z = 4$ the difference is slight. Dwarf 3 has a similar behaviour to dwarf 2, except its ratio drops relative to the no feedback simulation, eventually lying a factor of a few lower.

Dwarf 4 is the only case where there is a dramatic suppression of star formation by feedback. This object has a major merger (with a ratio $\sim 1.5$) around $z = 5.5$, so we treat the two progenitor haloes separately prior to their merger. Without feedback, both progenitor haloes have similar stellar to halo mass ratios to the other dwarfs, although they are individually of lower mass. With the inclusion of SN feedback, the ratio is dropped by approximately an order of magnitude at $z = 10$ and this offset increases with time. By $z = 4$, the progenitor haloes have merged according to the halo finder, although the central galaxies of the progenitors have not yet merged. For consistency, we now calculate the stellar to halo mass ratio for the final halo by considering the stellar mass of both of these galaxies. The ratio is now a factor of $\sim 100$ lower than the simulation without feedback and is close to the abundance matching relations (bearing in mind their uncertainties at this mass). The reason for the increased effectiveness of the feedback in this case would appear to be that this object has evolved for most of history as two independent systems that are less massive at a given redshift than the other simulated dwarfs, the shallower potential well increasing the relative efficiency of the SNe to clear gas. This suggests that haloes in this mass range are very sensitive to the manner of their assembly.

Having discussed the integrated efficiency of star formation, we now consider the star formation histories of our dwarfs shown in Fig. 3.4. For dwarf 4, we plot results for both central galaxies as in Fig. 3.3. For dwarf 1, the SFRs are essentially the same in the runs with and without SNe. Star formation starts at around $z = 11.5$
Fig. 3.4 Star formation rates as a function of redshift for the central galaxies. For dwarf 4, the central galaxies of the two almost equal mass progenitor haloes are shown, including when the two subhaloes are present in the same halo after the merger (see main text for more details); the second of the two haloes is plotted with a dashed line. Mechanical feedback in general leads to more bursty star formation rates.

and rapidly climbs to $2 \times 10^{-1}$ M$_\odot$ yr$^{-1}$ by $z = 10$. This rapid rise in star formation coincides with a merger at $z \sim 11$. The SFRs remain around this level until $z = 4$ in the no feedback run, apart from a merger-driven increase at $z \approx 5.5$. The results of this merger are apparent in the highly disrupted gas and disc structure visible in Fig. 3.2. With SNe, the brief increase in SFR is arrested by the feedback and dropped well below the no feedback rates. This burst of feedback is responsible for the more diffuse gas apparent in Fig. 3.2. The ability of the feedback to be effective during the later merger but not during the first merger is due to the amount of gas available. The subhalo gas fraction (relative to the total mass) at $z = 5.5$ is approximately a quarter of that at $z = 11.5$.

In contrast to dwarf 1, SNe are able to suppress the SFR significantly in dwarf 2 above $z = 6$. Without feedback, the SFR rises in a similar manner to dwarf 1, although not as rapidly. However, SNe are able to restrict star formation to
a brief burst at around $z = 11.5$ and another at $z = 8$. It would appear that the calmer environment (i.e. no major merger), relative to dwarf 1, at the onset of star formation allows the SNe to be effective. Dwarf 2 experiences a gas rich merger around $z = 6$ that leads to a large spike in SFR in both no feedback and feedback runs and the rapid build up of gas overwhelming the feedback. Following this event, the SFR remains high in both runs, leading to the similar (high) stellar mass to halo mass ratio at $z = 4$. A burst of efficient feedback around $z = 4.5$ leads to a slight drop in SFR relative to the no feedback simulations, the results of which can be seen in the gas morphology in Fig. 3.2.

In dwarf 3, without feedback, the SFR rises slowly from $z = 14$, before becoming reasonably steady at a few $10^{-1}$ $M_\odot$ yr$^{-1}$ from $z = 10$ onwards. This dwarf experiences no mergers of consequence, growing more slowly than dwarfs 1 and 2, probably as a result of being in a less dense environment. Once SNe are included, SNe are able to suppress star formation, but only following extended bursts of high SFRs. The feedback episodes are able to remove gas from the centre of the halo (giving rise to the morphology seen in Fig. 3.2) and the lower final stellar mass. However, a sufficiently large mass of stars is formed in the bursts such that the galaxy still lies several orders of magnitude above the (extrapolated) abundance matching relations.

As mentioned previously, despite ending up with a $z = 4$ halo mass similar to the other dwarfs simulated, dwarf 4 spends most of its history as two lower mass systems prior to a late major merger. Correspondingly, in the runs without feedback, the progenitors have lower SFRs than the other dwarfs, although this results in similar stellar to halo mass ratios (see Fig. 3.3). Peaking at $5 \times 10^{-2}$ $M_\odot$ yr$^{-1}$ by $z = 9$, the SFR of both galaxies evolves in a similar fashion. There is a slight drop in SFR after $z = 9$. The two haloes merge around $z = 5.5$ leading to a rapid increase in star formation. Like dwarfs 2 and 3, with the addition of feedback, the initial onset of star formation is limited to a short burst. However, the system is even more efficiently cleared of gas, resulting in a complete lack of star formation until the merger occurs. Unlike the no feedback case, this merger is relatively dry so the merger-triggered star formation burst is severely curtailed.

Dwarf 5 starts forming stars at $z \approx 15$, rising to high SFR after $z = 10$. A large amount of variability can be seen, mainly corresponding to mergers. Feedback has very little impact on the SFR in an averaged sense, although it impacts the gas near the very centre of the halo enough to cause variations relative to the no feedback simulation.

Fig. 3.5 shows circular velocity profiles, total density profiles and the ratio of
Fig. 3.5 Radial profiles at $z = 4$ for the various simulations, solid and dashed lines indicating runs with and without SNe, respectively. Vertical dotted lines denote the virial radius. Dwarf 5 without feedback is not shown as it was halted prior to this redshift, however its $z = 5.5$ profiles are consistent with the results from the other dwarfs. Top left: circular velocity profiles. While SN feedback systematically reduces the peak of circular velocity profiles, this reduction is only moderate (with the exception of dwarf 4). Top right: total (i.e. dark matter, gas and stars) density profiles. Dotted lines show profiles from collisionless dark matter only simulations. Bottom: the ratio of dark matter density in the simulations with baryonic physics to the collisionless simulations (renormalized by the cosmic dark matter fraction).

dark matter density in simulations with baryonic physics compared to collisionless (i.e. dark matter only) simulations at $z = 4$. It can be seen that on the whole, the simulations give rise to extremely concentrated mass distributions. The circular velocity profiles are strongly peaked at very small radii (10s of parsecs), in some cases $> 100 \text{ km s}^{-1}$. The inclusion of SNe reduces the magnitudes of the peaks by a factor of a few. Dwarf 4, which has managed to significantly suppress star formation (as seen in Figs. 3.3 and 3.4), is unique in preventing a peaked circular velocity profile. Instead, a gently rising profile reaches its peak value $\sim 30 \text{ km s}^{-1}$
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Fig. 3.6 Projected stellar half-mass radius vs. stellar mass for the central galaxies at $z = 10, 8, 6$ and $4$. Once again, in the case of dwarf 4, the central galaxies of the two progenitor haloes are shown. The points are calculated as the mean over a sample of 500 random viewing angles, with the error bars marking one standard deviation (open symbols are for no feedback runs, while filled symbols are for simulations with SNe). Horizontal dotted lines indicate the gravitational softening length at a given redshift (at $z = 4$, the blue dotted line corresponds to dwarfs 3, 4 and 5). Also plotted are observations for local dwarfs (McConnachie 2012; Koposov et al. 2015) for comparison, although a comparison of these $z = 0$ objects with our $z = 4$ galaxies should be treated with caution. Most objects form the majority of their stars in a dense central region limited only by the softening length (dwarf 3 no feedback is an outlier, see text). SNe have little impact, except in dwarf 4.

near the virial radius where it converges with the no feedback profile.

The centrally concentrated mass distribution that gives rise to these strongly peaked circular velocity profiles can be seen in the central panel of Fig. 3.5 in the form of radial profiles of total density (i.e. dark matter, gas and stars). Also plotted are profiles from collisionless simulations (dotted lines). These latter profiles are
fit well by NFW profiles (Navarro et al. 1997). The introduction of baryons leads to a strong peak of gas and stars with 0.1 kpc, which is overdense relative to the collisionless simulations by a factor of 100 in the centre. While the baryonic mass is dominant in this region, it can be seen in the rightmost panel of Fig. 3.5 that dark matter density has also been enhanced by a factor of $\sim 10$. Here, we plot the ratio of the dark matter density to the density from the collisionless simulations (renormalized by the cosmic dark matter fraction). The central concentration of baryons has led to contraction of the dark matter. Only in dwarf 4 has the feedback managed to expel sufficient baryons to prevent this central overdensity, its total and dark matter density profiles lying marginally under the collisionless case.

Fig. 3.6 shows the 2D projected stellar half-mass radius, $R_{1/2}$, as a function of stellar mass for the various galaxies at $z = 10, 8, 6$ and 4. We make this measurement from 500 randomly distributed viewing angles. The mean of the sample is plotted, error bars indicating the $1\sigma$ limit of the distribution. We mark with horizontal dotted lines the gravitational softening lengths. For reference, we also plot observations of local dwarfs (McConnachie 2012; Koposov et al. 2015), although the comparison of these $z = 0$ objects with our $z = 4$ dwarfs should be taken with some caution. The majority of our galaxies have extremely compact stellar distributions. While the projections in Fig. 3.2 show extended discs on the scale of hundreds of parsecs, most of the stellar mass is contained within a few tens of parsecs. In fact, the stellar half-mass radius is very close to the gravitational softening length, indicating that that the objects have undergone catastrophic collapse halted only by our limited resolution. The two component subhaloes of dwarf 4 remain less concentrated with the inclusion of SNe, lying at $z = 4$ within a factor of a few of the $z = 0$ observations at (dwarfs 2 and 3 also have larger $R_{1/2}$ at $z = 10$ before the failure of the SNe at later times).

Fig. 3.7 shows the location of our objects on a Kennicutt-Schmidt plot (SFR surface density as a function of gas surface density) between $z = 12 - 4$. We make these global measurements by taking the face-on ‘disc’ projection defined by the total angular momentum vector of the gas within twice the 3D stellar half-mass radius (although it should be noted that not all of our galaxies produce discs). For a given projection of the galaxy, we find the 2D radius containing 90% of the total mass of stars.

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3The results are well converged with number of samples above 500.

4Dwarf 3 without feedback at $z = 4$ is somewhat of an outlier. The lack of any disruption from mergers has allowed a more extended stellar distribution. As the majority of these stars are formed between $z = 6 - 4$, this leads to a sudden increase in $R_{1/2}$ by $z = 4$. The simulation with SNe has a significantly smaller $R_{1/2}$, but this is mainly because it has a proportionally lower mass of stars.
Fig. 3.7 A Kennicutt-Schmidt plot, showing SFR surface density as a function of gas surface density for our simulated dwarfs between $z = 12 - 4$ ($z = 12 - 5.5$ for dwarf 5 with no feedback). We only plot the most massive of the two dwarf 4 subhaloes for clarity, but the secondary subhalo exhibits similar behaviour. These are global measurements, taken within a radius containing 90% of the total SFR, projecting down the gas angular momentum vector (open symbols are for no feedback runs, while filled symbols are for simulations with SNe). Also shown are observations, both global (Kennicutt 1998; Wyder et al. 2009) and spatially resolved (Bigiel et al. 2008). We also plot the power law fit of Kennicutt (1998) to the data of that work. Most of our simulated galaxies have high gas surface densities and SFR surface densities. A few galaxies experience strong bursts of feedback which drive them well beyond the boundaries of the plot as they are quenched, the few low surface density points representing transitions.

We then compute SFR surface density and mass surface density from the gas within this radius. Fig. 3.7 also shows global (Kennicutt 1998; Wyder et al. 2009) and spatially resolved (Bigiel et al. 2008) observations. Due to the extremely compact nature of most of our galaxies, the majority of our simulations appear in the same region of the Kennicutt-Schmidt plot as starburst galaxies. There is a trend for our simulations with feedback to produce galaxies with slightly lower SFR and mass surface densities than simulations without feedback. When galaxies

\footnote{The measurements are relatively insensitive to the exact fraction adopted, the points being shifted up and down the Kennicutt-Schmidt relation slightly.}
experience an efficient burst of feedback (dwarfs 2, 3 and 4; see Fig. 3.4) they move towards the lower end of the relation. However, because these bursts are very strong and tend to completely disrupt the star forming gas, we do not see a steady state at low surface densities, but the measurements in this quenched phase lie well beyond the boundaries of the plot. For example, dwarf 4 with feedback only appears on the plot at \( z = 11.5, 11 \) and 4.5 because it effectively has no star formation at other times (we do not plot the secondary subhalo of dwarf 4).

Fig. 3.8 shows kinematic information of our simulated galaxies as a function of stellar mass at \( z = 4 \) as compared to measurements of local dwarfs from Wheeler et al. (2017). The top left panel shows the rotational velocity. We take here the peak value of the stellar rotation curve, having first transformed into the ‘disc’ plane of the galaxy by aligning with the total angular momentum vector of the stars. It should be noted that the kinematics from the simulations should be treated with caution given that the size of the systems approaches the gravitational softening length in those cases in which catastrophic collapse has occurred. The rotational velocities are well in excess of the observations, but not unexpected given the highly peaked circular velocity profiles (see Fig. 3.5). There is a trend for the simulations with SN feedback to produce higher rotational velocity systems (except for dwarf 4). With SNe, however, the two subhaloes of dwarf 4 show no evidence of rotation and are therefore consistent with the observations at that mass which demonstrate little or no rotation.

The top right panel of Fig. 3.8 shows the 1D velocity dispersion, \( \sigma \), for our systems. We measure the 3D velocity dispersion within a sphere whose radius corresponds to the peak rotational velocity\(^6\), then obtain the 1D value by dividing by \( \sqrt{3} \). In the case of dwarf 4 with SNe (which shows no rotation) we use the stellar half-mass radii. Again, most simulations lie significantly above the local observations, a consequence of the highly compact systems (Fig. 3.6 demonstrates how much more extended observed galaxies in this mass range are). There is a steep relation of increasing \( \sigma \) with increasing stellar mass. The two subhaloes of dwarf 4 with SNe lie close to the observations, with velocity dispersions of \( \sim 5 \text{ km s}^{-1} \).

Examining the ratio of the rotational velocity to the velocity dispersion provides a measure of the rotational support of the system. This is shown in the bottom panel of Fig. 3.8. Most of our systems are rotationally supported, in contrast with the observations which prefer rotation to be subdominant (although there are a few outliers and the uncertainties are large in some observations), with the caveat that

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\(^6\)Taking other reasonable radii, such as one or two times the stellar half-mass radius yields the same results within 1 km s\(^{-1}\); these radii are all comparable.
Fig. 3.8 Stellar kinematic properties of the simulated dwarf galaxies at $z = 4$. 

Top left: the peak rotational velocity, having aligned the system in the ‘disc’ plane relative to the total stellar angular momentum vector. Top right: the 1D velocity dispersion within the peak rotational velocity radius (or in the case of dwarf 4 with SNe, within stellar half-mass radius, see text for details). Bottom: the ratio of rotational velocity to velocity dispersion, a measure of rotational support. Also plotted are measurements of observations from Wheeler et al. (2017) of dwarfs in the Local Group.

we are comparing our $z = 4$ objects with local observations. Only dwarf 4 with SNe is consistent with observations producing a dispersion dominated system.

Because dwarfs 1, 2, 3 and 5 produce such large masses of stars in a confined region, the resulting metal enrichment of the surrounding region is necessarily extremely high. Fig. 3.9 shows radial stellar metallicity profiles at $z = 4$. Without SNe, the central tens of parsecs (which contain most of the stellar mass) are dominated by a stellar population of $\sim 2 Z_\odot$. The metallicity drops rapidly through the disc region (on the order of 100 pc, see also Fig. 3.2) to reach a metallicity ranging between 0.2 $Z_\odot$ (dwarf 4) and 0.6 $Z_\odot$ (dwarf 2) in the stellar ‘halo’. The metallicity gradient is then flat until the edge of the stellar distribution, after which the pro-
Fig. 3.9 Radially (mass-weighted) averaged stellar metallicity profiles at $z = 4$. Outside of a few kpc, the profiles become very noisy, in some cases because of substructures. For dwarf 4, the profiles are centred on the most massive subhalo from the no feedback simulations. While in dwarfs 1, 2, 3 and 5 inefficient SN feedback leads to over-enrichment, in the case of dwarf 4 this is reduced by two orders of magnitude to more reasonable values.

Files are noisy due to the low stellar density and the presence of other subhaloes. With the exception of dwarf 4, the addition of SNe makes very little difference to the stellar metallicities, which is to be expected given the inefficiency of feedback in these systems (although, in dwarf 3, the stellar halo is curtailed at a smaller radius). For dwarf 4, SNe reduce the central metallicities by 2 orders of magnitude to $0.02 - 0.03 Z_{\odot}$. The resulting metallicity gradient is flat through the entire stellar distribution, out to $\sim 1$ kpc (the second subhalo appears in this radial profile at larger radii, as can also be seen in the no feedback profile). While the lower stellar metallicities are partially due to the lower overall stellar mass formed relative to the no feedback simulation, the ability of the SNe to expel metals from the centre of the halo is also key.

Fig. 3.10 shows mass-weighted gas metallicity projections of dwarf 4, without and with the inclusion of feedback. In the absence of feedback, metals remain where they have been deposited by the star particles, leading to high concentrations around the subhaloes. This can be seen in the projection, where the majority of gas
(both inside and outside of the virial radius) remains at the metallicity floor of the initial conditions, $10^{-4} Z_\odot$. Very small patches of super-solar metallicity gas can be seen around the subhaloes, while some metal enriched gas has been stripped during the merger, leaving trails. With the inclusion of feedback, gas of a few $10^{-2} Z_\odot$ is widely distributed inside and outside of the virial radius. In the other dwarfs, the inclusion of feedback also allows metals to leave the halo ($\sim 10^{-1} Z_\odot$ at the virial radius), but approximately 3 orders of magnitude more stellar mass has been created to achieve this i.e. the SNe are $\sim 100$ times less efficient at ejecting metals. We reported a similar phenomenon in isolated simulations in Chapter 2, where inefficient SNe lead to slow moving, highly metal enriched outflows simply due to the number of SNe occurring. Nonetheless, dwarf 4 demonstrates that it is possible for dwarfs to efficiently enrich the CGM.

Fig. 3.11 shows gas mass outflow and inflow rates across $0.25 R_{\text{vir}}$, $0.5 R_{\text{vir}}$ and $R_{\text{vir}}$ as a function of redshift for dwarf 4 with and without SN feedback. SFRs are also plotted for comparison. The outflow rates are calculated as:

$$\dot{M}_{\text{out}} = \frac{\sum_i m_i v_{\text{out},i}}{\Delta r},$$

(3.3)
where the sum is over all gas cells within a shell of thickness $\Delta r = 50$ pc centred on the target radius that have a positive radial outflow velocity, $v_{\text{out}}$. The inflow rates are calculated in the same manner, but for all cells that have a negative radial velocity. In the absence of SN feedback, outflow rates mostly remain well below inflow rates, with outflows only arising from the motion of substructures and mergers. For example, the dramatic increase in outflow rates just before $z = 4$ is due to the motion of the merging subhalo within the primary halo. With SN feedback, after the bursts of star formation, outflow rates increase dramatically while inflow is suppressed.

Outflows are often characterised in terms of mass loading factor, i.e. the ratio of outflow rate to SFR. The outflows across the three radii are offset from the peak of the SFR due to the time difference between star formation and SNe exploding as well as the travel time of the outflow, so an instantaneous mass loading factor is not a useful quantity. However, comparing the peak SFRs and outflow rates yields a mass loading factor of approximately 90, 60 and 30 across $0.25R_{\text{vir}}$, $0.5R_{\text{vir}}$ and $R_{\text{vir}}$, respectively. Following the burst of star formation at $z = 11$, inflow across $0.25R_{\text{vir}}$ is essentially halted until after $z = 10$. The inflow rates remain a factor of $\sim 5$ below the corresponding no feedback simulation rates until the merger begins at $z \approx 5.5$. At this point, it appears that the UV background is hindering the ability
of gas to condense into the centre of the halo. The second burst of star formation after $z = 5$ also produces a brief outflow, preventing further star formation. None of the other dwarfs simulated are able to produce strong outflows. Dwarfs 2 and 3 have very brief outflows after bursts of star formation and subsequent efficient feedback, but they have mass loading factors $< 2$ and barely suppress inflow rates except in the very centre of the halo.

3.4 Discussion

3.4.1 Why is SN feedback inefficient?

In all of our simulated galaxies except for dwarf 4, SN feedback is unable to prevent the catastrophic collapse of gas and resulting runaway star formation. The reason for this inefficiency appears to be that most SNe occur in very dense gas. This can be seen in Fig. 3.12 which shows the distribution of gas density in which SN explode for all SNe above a given redshift for dwarfs 1 and 4, with and without feedback. Comparing the PDFs for dwarf 1 at $z = 11$ for the no feedback and feedback simulations, both peak at a high density of $\approx 10^3$ cm$^{-3}$. With feedback, there is a slight tail to low density, indicating that the feedback has been able to clear some gas. A short time later at $z = 10$, the peak of the distribution is at $\approx 10^4$ cm$^{-3}$. By contrast, while without feedback dwarf 4 has a similar PDF to dwarf 1, once SNe are included the peak of the distribution is at $\sim 3$ cm$^{-3}$ for all redshifts. The decisive criterion determining the success and failure of the feedback is whether it is able to clear the dense gas immediately, otherwise the SNe will be overwhelmed.

Dwarf 1 fails the criterion immediately, as does dwarf 5. Dwarf 2 succeeds twice but is overwhelmed by a sudden increase of gas during a wet merger at $z \sim 6$. Dwarf 3 is partially successful, but the bursts reach too high a SFR before the system is quenched, so the net reduction in stellar mass is too low. Dwarf 4 is unique amongst our simulations in being completely successful, mainly due to its merger history. As described in the previous section, while the final halo is comparable in mass to dwarf 1, it is formed from a major merger (with a mass ratio $\sim 1.5$) late in its history ($z \approx 5.5$). This means that it spends most of its evolution as two smaller haloes. This makes it easier to clear gas for two reasons. Firstly, there is a shallower potential well to fight against. Secondly, the inflow of gas onto the haloes is reduced relative to dwarf 1 (even without feedback, the SFRs for the two progenitor haloes are lower than for dwarf 1; see Fig. 3.4).
Fig. 3.12 Distribution of the densities of gas in which SNe occur. Dwarf 1 and 4 are compared, with and without feedback. The redshift evolution of the PDFs are shown (cumulatively). For numerical reasons, these PDFs are for all SNe that occur in the high-resolution region, rather than being tied explicitly to the host halo of a given dwarf. However, the vast majority of SNe occur in the host halo, so these PDFs are representative. Most SNe occur in gas with a density of approximately $10^4$ cm$^{-3}$ for dwarf 1 (with and without feedback) and dwarf 4 without feedback. However, with the inclusion of SN feedback in dwarf 4, the mean density drops by three orders of magnitude.

It may appear at first glance that the inefficiency of SNe in dense gas is a result of shortcomings in our method of feedback injection i.e. numerical overcooling. However, our mechanical scheme is designed to mitigate the effects of under-resolved SN remnants by injecting the correct momentum relative to the stage of their evolution that can be resolved. Full details can be found in Chapter 2 where we also demonstrate using isolated simulations that this scheme is numerically robust (see also the following section where we discuss convergence with resolution). For extremely dense gas, at most tractable resolutions, the SN remnant will remain entirely unresolved so our scheme will inject the final momentum achieved during the Sedov-Taylor phase. We make use of a fitting function to high resolution simulations of individual SNe (see Blondin et al. 1998; Thornton et al. 1998; Geen et al.
2015; Kim & Ostriker 2015; Martizzi et al. 2015; Kimm et al. 2015),

\[ p_{\text{fin}} = 3 \times 10^5 \text{ km s}^{-1} \text{ M}_{\odot} E_{51}^{16/17} n_{\text{SN}}^{-2/17} Z_{\text{SN}}^{-0.14} , \]  

(3.4)

where \( E_{51} = (E_{\text{SN}}/10^{51} \text{ ergs}) \) is the energy of the SN (for our individually time-resolved SNe, \( E_{51} \equiv 1 \)), while \( n_{\text{SN}} = (n_{\text{H}}/\text{cm}^{-3}) \) and \( Z_{\text{SN}} = \text{MAX} (Z/Z_{\odot}, 0.01) \) are the hydrogen number density and metallicity of the ambient gas, respectively.

It can therefore be seen by comparing the peaks of the density PDF for SN sites that in dwarf 1 the momentum budget per SN is reduced to \( \sim 0.39 \) of that in dwarf 4. Additionally, if metals are not cleared efficiently this will also impact the efficiency. Given that the typical gas metallicity in the centre of dwarf 1 is approximately a factor of 100 higher than in dwarf 4, this reduces the momentum budget again by half, meaning that in total only 20% of the momentum budget per SN is available relative to dwarf 4. In addition to impacting the small scale evolution of the SN remnant, the build up of a central concentration of dense gas will make it more difficult for the momentum injection from SNe to clear material from the galaxy because the mass of material that must be swept up in order for an outflow to escape becomes proportionally higher. These two factors lead to a state of runaway star formation if at any point the feedback is unable to prevent the build up of dense gas, particularly if inflow rates increase suddenly (e.g. due to mergers, be they major or minor).

### 3.4.2 The impact of the choice of parameters on our results

Having discussed the reasons why SN feedback is inefficient in our fiducial simulations, we now explore the degree to which our results are generally applicable as opposed to being dependent on our choice of parameters. Fig. 3.13 shows the SFR as a function of redshift and the stellar mass to halo mass ratio as a function of halo mass (at \( z = 10, 8, 6 \) and 4) for our fiducial simulations of dwarf 1 as well as 8 resimulations in which we alter various parameters of our models.

Increasing the number of resolution elements in the zoom-in region by a factor of \( 2^4 \), giving a mass resolution of 191 M\(_{\odot}\) and 35.9 M\(_{\odot}\) for dark matter particles and gas cells, respectively, has very little impact on the results. While the SFR shows slightly more variation than the fiducial resolution simulation and there is a suppression of star formation briefly between \( z = 5.5 - 5 \), even with the increased resolution the feedback is unable to prevent runaway star formation beginning at early times. This leads to a \( z = 4 \) stellar mass that only differs from the fiducial simulation by a factor of 1.2. The convergence properties of our mechanical feed-
Fig. 3.13 SFR as a function of redshift (top) and integrated star formation efficiency (bottom) for simulations of dwarf 1 with a variety of alternative parameters. The results are split into two columns for clarity. No feedback (black) and fiducial SNe (red) simulations are repeated in both panels for reference. The additional simulations are as follows: left: we increase the mass resolution in the zoom-in region by a factor of 8 to 191 M⊙ and 35.9 M⊙ for dark matter particles and gas cells, respectively (green), the Faucher-Giguère et al. (2009) UV background is turned on from our first available tabulated redshift of $z = 11.7$ (blue), we impose a cap of 100 cm$^{-3}$ on density that is used to determine the maximum momentum that can be injected for a SN (see eq. 3.4) (yellow); right: the pressure floor is turned off (cyan), the star formation density threshold is increased by a factor of 10 to 100 cm$^{-3}$ (pink), the star formation efficiency is increased by a factor of 10 to 15% (brown), the pressure floor is turned off and the star formation efficiency is set to 100% (orange). Abundance matching relations at $z = 4$ (Behroozi et al. 2013; Moster et al. 2018) are shown, although they are extrapolated into this mass range.
back scheme as a function of resolution were demonstrated in a non-cosmological isolated galaxy setup in Chapter 2, albeit in a regime where the feedback was capable of regulating star formation.

We can see that our choice to delay turning on the Faucher-Giguère et al. (2009) UV background until $z = 9$ is similar to switching it from $z = 11.7$, apart from a slight reduction in SFRs between $z = 10 - 9$. This shows that the assumed UV background is unable to prevent the catastrophic build up of gas at $z = 10$. We note, however, that this conclusion rests on the approximation of a homogeneous UV background as opposed to local radiation fields. We have neglected photoionisation from the stars formed in the galaxies themselves. This may be able to prevent the build up of dense gas in star forming regions (see e.g. Vázquez-Semadeni et al. 2010; Walch et al. 2012; Dale et al. 2014; Sales et al. 2014). In addition, dwarfs that are in crowded regions or are satellites of larger galaxies may be bathed in ionising radiation from nearby external sources, assuming that those galaxies are able to clear/ionise sufficient local gas to achieve a high enough escape fraction for UV photons. The failure of the UV background to quench our dwarfs is not inconsistent with other works that indicate the existence of a $z = 0$ threshold mass of a few $10^9 M_\odot$ below which UV background is effective (see e.g. Okamoto et al. 2008; Shen et al. 2014; Sawala et al. 2014; Wheeler et al. 2015; Fitts et al. 2017) as our dwarfs will have $z = 0$ virial masses in excess of $10^{10} M_\odot$.

We further carry out a simulation in which we modify the equation used to calculate the final momentum of a SN remnant after the Sedov-Taylor phase (eq. 3.4) such that the dependence on ambient gas density is capped at 100 cm$^{-3}$. This is a crude approximation to the idea that local stellar feedback may have prevented surrounding gas reaching high density prior to the first SN occurring. Imposing this density cap increases the momentum budget per SN by a factor of 1.7 relative to SN occurring in gas with $10^4$ cm$^{-3}$ (as in Fig. 3.12). Of course, the gas itself is still at high density, so continues to present an obstacle to efficient clearing of material from a hydrodynamical standpoint. Nonetheless, with this caveat in mind and very moderate increase in the momentum budget, this simulation results in a factor of 3 lower stellar mass at $z = 4$. This hints that the need to regulate local gas density is important, but also demonstrates that such a simple modification to the subgrid scheme is not sufficient to obtain realistic galaxy properties.

The use of pressure floors to prevent artificial fragmentation is a subject of some debate in the literature. We discussed the impact of adopting such a technique in some detail in Chapter 2, so we refrain from an in-depth discussion here. Nonetheless, we tested the impact on our zoom-in simulations by resimulating dwarf 1
without a pressure floor. As can be seen from Fig. 3.13, this has a negligible impact on our results. The lack of a floor seems to produce slightly more clustered SNe, leading to a reduction in SFR by a factor of a few from $z \sim 10 - 7$. However, in general the SFR is similar to the fiducial simulation and the $z = 4$ stellar mass is the same within 4%.

Increasing the star formation threshold by an order of magnitude to $100 \, \text{cm}^{-3}$ also produces more clustered SNe at early times, allowing the feedback to quench star formation at $z = 10$. This leads to a reduction in stellar mass relative to the fiducial case by a factor of a few at $z = 8$. However, this is still not enough to prevent the build up of dense gas at later times. From $z = 7$ onwards, the SFR is similar to the fiducial case, leading to a reduction of the $z = 4$ stellar mass by only 1.3.

Increasing the star formation efficiency, $\epsilon_{\text{SF}}$ by a factor of 10 to 15% leads to significantly different behaviour. The SFR rises faster and strong clustering of SNe leads to efficient launching of outflows and the suppression of star formation. Star formation proceeds in short bursts for the entire duration of the simulation. Despite this, the $z = 4$ stellar mass is only reduced by slightly over an order of magnitude, leaving it over 2 orders of magnitude above the (extrapolated) abundance matching relations and an order of magnitude larger than dwarf 4 with the fiducial star formation parameters. Failing to match the abundance matching relations at this redshift is not necessarily a failure in and of itself because of the uncertainties involved at this mass range. If the halo was to subsequently form very little stellar mass it may obtain a stellar mass to halo mass ratio consistent with observations at lower redshift (where the stellar mass-halo mass relation is more constrained). However, given that the simulation is still forming stars towards $z = 4$ and that the mass ratio is similar at $z = 6$ and $z = 4$ it remains unclear that this would be the case.

Finally, we try an extreme choice of parameters in an attempt to reduce the stellar mass further. We turn off the pressure floor and use $\epsilon_{\text{SF}} = 100\%$. This leads to extremely rapid star formation and a concentrated burst of SN feedback that is able to completely quench the galaxy, expelling most of the gas. Star formation does not resume by $z = 4$. The result is a reduction in $z = 4$ stellar mass by almost 2 orders of magnitude. While this is still too high relative to the extrapolated abundance matching relations, it is possible that this galaxy would move onto the relation at lower redshift. While this may be seen as a successful solution as the choice of $\epsilon_{\text{SF}}$ may still be within the physically plausible range (see paragraph below), a more cautious interpretation would indicate that, given that we need to
push our star formation model to its extremes in order to be successful, we are likely neglecting some other important physical processes that would alleviate the need for very high values of $\epsilon_{SF}$ in the first place.

The reason for the increase in feedback efficiency as a result of increasing $\epsilon_{SF}$ is twofold. Firstly, it leads to more clustered SNe that are able to work together to drive outflows. Secondly, it avoids the issue of building up high density gas by efficiently converting gas into stars before the problem arises. Care, however, must be taken when using such high values of the efficiency that this does not represent an unphysical removal of gas. As the gas consumption time is then effectively the free-fall time, above $100 \text{ cm}^{-3}$ this becomes comparable to the time before the first SNe explode, meaning that most, if not all, of the local gas will have been converted into stars, significantly dropping local density for subsequent SN events. If the internal structure of star-forming regions is well resolved this may not be particularly problematic because the hydrodynamics should correctly follow the fragmentation of the region without recourse to 'fudge factors'. However, if the region is unresolved, using an efficiency of 100% will quickly convert the entire mass of the region into stars, which is likely unphysical. To avoid this effect, high efficiencies could perhaps be used in combination with more complex selection criteria for star forming gas (for example using virial parameters, checking whether the region is Jeans unstable (Hopkins et al. 2014)), but only if the GMC is at least partially resolved. Alternatively, variable star formation efficiencies that depend on gas properties (e.g. turbulent state, virial parameter (Federrath & Klessen 2012; Kimm et al. 2017)) may be able to bridge the gap between fully resolved and partially resolved simulations of star forming regions.

It is worth reemphasizing that regardless of the star formation criteria, there is a large body of theoretical and observational work indicating that other sources of stellar feedback must be operating prior to the first SN, such as stellar winds, photoelectric heating and photoionization from young stars. These processes may have a significant impact on local gas, not only affecting its density and temperature structure, but also the level of turbulent support. Given that we have demonstrated a tendency for dense gas to build up and overwhelm SN feedback in our $z = 4$ dwarfs (and that this effect is physically realistic, rather than just being a symptom of numerical overcooling), it may be the case that non-SN stellar feedback plays a more important role in the evolution of low mass haloes than is commonly assumed. This conclusion is consistent with the results found by the FIRE-1 project (Hopkins et al. 2014) in which the removal of other sources of stellar feedback in dwarfs led to SN feedback having almost no impact on stellar mass (though the effect
appears to be less severe in FIRE-2 (Hopkins et al. 2017)). Finally, we note that the efficiency of first (and subsequent) SN events may depend on the fraction of runaway SN and on alternative heating processes such as those provided by relativistically accelerated particles in the wake of SN explosions.

### 3.5 Conclusion

We have carried out very high resolution cosmological zoom-in simulations of five dwarf galaxies up to $z = 4$ with virial masses between $\sim 2 - 6 \times 10^9$ M$_\odot$. Our simulations adopt the mechanical SN feedback scheme introduced in Chapter 2 and a spatially constant, but time evolving UV background (Faucher-Giguère et al. 2009). The SN feedback is constructed to deliver the correct momentum to the surrounding ISM corresponding to the stage of the SN remnant evolution. We found that this model leads to self-regulated star formation rates, realistic galaxy kinematics and gas content thanks to the occurrence of multiphase, mass-loaded outflows in isolated dwarf simulations. The aim of the present work is to determine whether the exact same model of SN feedback results in realistic dwarf properties once the full cosmological formation is incorporated self-consistently. We find that:

- Without the inclusion of SN feedback, we produce dwarfs that have over 3 orders of magnitude too much stellar mass relative to (extrapolated) abundance matching predictions. Their stellar and gas metallicities are in excess of solar abundances. The dwarfs undergo a catastrophic collapse to the resolution limit, resulting in extremely dense systems with strongly peaked circular velocity curves. Dark matter density in the centre of the halo is enhanced relative to a collisionless simulation by approximately an order of magnitude.

- In general, while the inclusion of SN feedback induces more bursty SFR rates and affects dwarf morphologies, it has insufficient impact on the total stellar mass formed. In the majority of our systems, the build up of dense gas (often following a wet merger) renders the SNe too inefficient to expel gas from the galaxy and suppress star formation. We emphasise that, because our scheme injects the correct amount of momentum per SN, this effect is not an example of classical numerical overcooling but rather a physical suppression of SN efficiency. Most SNe explode in gas of density $10^4$ cm$^{-3}$ which limits the feedback momentum budget available. This suggests that some other mechanism(s) must be invoked (e.g. other sources of stellar feedback) that can prevent gas from collapsing to such high densities prior to SNe occurring.
The inclusion of runaway SN may also help alleviate this issue.

- We however find one exception to this scenario where we are able to produce a realistic dwarf relative to the extrapolations of abundance matching and various metrics of local analogs. Our dwarf 4 forms by a major merger relatively late in its history at \( z \approx 5.5 \). Consequently, it spends most of its evolution as two lower mass systems in which the SNe are able to expel gas and halt star formation before catastrophic collapse sets in. Their late major merger is therefore mostly dry and does not trigger more than a brief burst of star formation which is quickly suppressed by feedback. We note that while SN feedback is clearly efficient here, enriching the CGM to a few \( 10^{-2} Z_\odot \) with mass-loaded winds, no prominent dark matter core forms.

- We have carried out a variety of other simulations to test the applicability of our conclusions. We find that our results are not significantly impacted by increasing resolution, changing details of the (spatially uniform) UV background or removing the pressure floor. Our results are also insensitive to increasing the star formation density threshold by an order of magnitude. Arbitrarily increasing the star formation efficiency parameter by an order of magnitude to 15% leads to more bursty behaviour and reduced star formation, but still overshoots abundance matching relations by 2 orders of magnitude. Only by taking an extreme choice of parameters, using a star formation efficiency of 100% and removing the pressure floor, are we able to get close to the relation.

We have demonstrated that realistically modelled SN feedback is easily overwhelmed early on in the cosmological assembly of dwarfs by the build up of gas, despite the relatively shallow potential well. While this can potentially be dealt with by adopting a star formation prescription that leads to extremely concentrated SN feedback, it seems that other sources of stellar feedback may be required to modulate ISM densities prior to the first SNe exploding in order to preserve their efficiency.
4. The impact of small-scale turbulence on star formation rates and supernova feedback

4.1 Introduction

As we discussed in Chapter 1, giant molecular clouds (GMCs) are the sites of star formation. Comparisons of their estimated freefall times (typically on the order of a few Myr or lower) to their gas consumption time (the ratio of SFR density to gas density, on the order of a few Gyr) or to their observed lifetime (a few tens of Myr, comparable to the lifetime of OB stars) suggest that they are stabilised against gravitational collapse on the cloud scale. Observations indicate that the gas in GMCs contains supersonic, compressible turbulence (see e.g. Zuckerman & Palmer 1974; Zuckerman & Evans 1974; Larson 1981; Solomon et al. 1987; Falgarone et al. 1992; Ossenkopf & Mac Low 2002; Heyer & Brunt 2004; Heyer et al. 2006; Brunt et al. 2009; Roman-Duval et al. 2011; Schneider et al. 2011; Ginsburg et al. 2013; Rathborne et al. 2014, 2015). It has been suggested that the turbulence acts to stabilise the clouds on large scales but causes shocks on small scales that induce compression, leading to the observed structure of filaments and cores that are the initial conditions for star formation (see e.g. Elmegreen & Scalo 2004; Mac Low & Klessen 2004; McKee & Ostriker 2007; Schneider et al. 2012). It is then possible that turbulence acts to regulate the cloud scale star formation efficiency (see e.g. Federrath & Klessen 2012, and references therein). In addition to altering the star formation efficiency directly, this will also affect the efficiency of SN feedback in disrupting star forming regions by altering the clustering of SNe (in space and time) (see e.g. Yadav et al. 2017; Gentry et al. 2017; Kim et al. 2017) and the porosity of the local medium (see e.g. Kim & Ostriker 2015; Martizzi et al. 2015; Walch & Naab 2015; Kimm et al. 2015). Possible drivers of turbulence include expanding supernova remnants (SNRs), galactic spiral shocks, gravitational contraction, local feedback from young stellar objects, galactic rotation and magnetorotational instabilities (MRI) (see Chapter 1 for references). Observed turbulence appears to be seeded at the largest cloud scales and proceeds via turbulent cascade to smaller scales (see e.g. Heyer et al. 2006; Brunt et al. 2009).

A great deal of work has been carried out investigating the turbulent ISM with
the cloud-scale simulations (see references above and in Chapter 1) and in stratified box simulations representing patches of galactic discs (see e.g. Gatto et al. 2015; Martizzi et al. 2015; Simpson et al. 2016). Building on the work presented in Chapter 2 we wish to examine how the level of turbulence in dense gas clouds impacts global galaxy properties. For example, how does a turbulent ISM structure impact the fragmentation of the disc? Is there an impact on the clustering of SNe and does this affect global disc structure and the mass loading factors of outflows on large scales? Unfortunately, in the typical isolated galaxy setup (as used in Chapter 2) we miss many of the drivers of cloud scale turbulence either because we do not include the necessary physics (e.g. the MRI requires MHD), our simple setup does not provide the requisite environment (e.g. we do not typically produce strong spiral arms), we do not include the full range of stellar feedback or we cannot resolve the small-scale ISM physics well enough to naturally produce turbulence.

We therefore present a new method in which we are able to produce idealised turbulent ISM properties within isolated galaxy simulations by the selective artificial driving of turbulence. The motivation here is to ignore the precise details of the mechanisms by which turbulence is driven, instead including a desired level of turbulence in the simulation ‘by hand’ and focusing on the consequences for global galaxy properties. We inject turbulence into a very narrow band at the cloud-scale and are able to choose which gas receives the driving for any arbitrary criteria (e.g. according to a density criterion to identify star forming regions). We can adjust the ratio of solenoidal to compressive turbulence. It is also possible to vary the amplitude of driving and the ratio of compressive to solenoidal driving based on the location of the gas within the disc (e.g. observations indicate that turbulence in GMCs in the Milky Way central molecular zone is largely solenoidal, whereas in spiral arms it is largely compressive, see e.g. Federrath et al. 2016, and references therein). For the purpose of this thesis, we present a very simple demonstration of the technique using initial conditions from Chapter 2. We do not attempt to calibrate the level of driving to any particular observational criteria and we refrain from a detailed analysis of the results in anything more than a qualitative sense.

4.2 Methodology

4.2.1 Code setup, star formation and feedback

We adopt the same basic setup as the previous two chapters, making use of the moving-mesh code AREPO with some modifications. Again, we use radiative cooling from primordial species and metals lines as in Vogelsberger et al. (2013). We do not include a UV background. Unlike the previous chapters we do not include a
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pressure floor because we wish to examine the impact of turbulent driving on pressure support. We continue to make use of a simple Schmidt law for star formation,

$$\dot{\rho}_* = \epsilon_{\text{SF}} \rho / t_{\text{ff}},$$  \hspace{1cm} (4.1)

where $\rho$ is the gas density, $\epsilon_{\text{SF}} = 1.5\%$ is the efficiency and $t_{\text{ff}} = \sqrt{3\pi/32G\rho}$ is the free-fall time. However, we increase our star formation threshold density, $n_{\text{SF}}$, to $100\,\text{cm}^{-3}$ to naively identify star forming regions with typical GMC densities (note, however, that we do not model formation of molecular hydrogen). However, we demonstrated in Chapter 2 that this is unlikely to make a significant difference to the resulting star formation rates. These star formation rates are stochastically sampled to convert gas cells into star particles.

As in Chapters 2 and 3, we tabulate SN rates as a function of age and metallicity from Starburst99 (Leitherer et al. 1999) using a Kroupa (2002) IMF. SN events are then triggered by stochastically sampling these rates, making use of a timestep limiter to make sure that SNe are always individually time-resolved. We deposit feedback quantities into the local gas using a weighting scheme that guarantees isotropy (see Chapter 2). We use an ejecta mass of $10\,\text{M}_\odot$ per SN of which $2\,\text{M}_\odot$ are metals. In runs described as being ‘no feedback’, we return mass and metals, but do not inject energy or momentum (other than that required to conserve the momentum carried by the star particle). When SN feedback is included, we use the mechanical scheme as described in Chapter 2. This injects the correct momentum according to the phase of the SNR evolution resolved.

4.2.2 Turbulent driving

We make use of a stochastic acceleration field to drive turbulence. Our scheme is based on an existing implementation in Arepo by Bauer & Springel (2012) (see also Fiacconi et al. 2018), designed for use in idealised periodic boxes. The procedure for generating the acceleration field follows that used by Schmidt et al. (2006); Federrath et al. (2008); Federrath et al. (2009, 2010); Price & Federrath (2010). The field is generated in Fourier space and the phases of the modes are drawn from an Ornstein-Uhlenbeck process (Uhlenbeck & Ornstein 1930). Developed to model Brownian motion, the process is a modification of a continuous-time random walk (or Wiener process) such that there is a drift back towards a central location (i.e. it is mean reverting). The corresponding random sequence is given by
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\[ x_t = \exp\left(\frac{-\Delta t}{t_s}\right) x_{t-\Delta t} + \sigma \left(1 - \exp\left(\frac{-2\Delta t}{t_s}\right)\right)^{1/2} z_t, \quad (4.2) \]

where \( t_s \) is the correlation time, \( \sigma^2 \) is the desired asymptotic variance of the Ornstein-Uhlenbeck process and \( z_t \) is a Gaussian random variable drawn from a distribution with unit variance. The sequence has zero mean, \( \langle x_t \rangle = 0 \), and time correlation \( \langle x_t x_{t+\Delta t} \rangle = \sigma^2 \exp\left(-\Delta t/t_s\right) \). This means that the time evolution of the acceleration field contains stochastic fluctuations but is smoothly varying over the correlation time. The Gaussian random variable \( z_t \) is generated by means of a Box-Muller transform of two independent random numbers \( U_1 \) and \( U_2 \) drawn from a uniform distribution in the interval \([0, 1]\):

\[ z_t = \sqrt{-2 \ln U_1 \cos(2\pi U_2)}. \quad (4.3) \]

We draw 6 independent \( x_t \) per mode with wave vector \( k \), corresponding to the three Cartesian components of the Fourier sine and cosine transforms, which we now express as the vectors \( x_s(k) \) and \( x_c(k) \) (we drop the subscript \( t \)), respectively. By means of a Helmholtz decomposition, we can express the acceleration field in terms of a combination of divergence-free (solenoidal) and curl-free (compressive) components. We can then alter the relative amplitude of each component to produce a field with a chosen mixture of compressive and solenoidal acceleration. In Fourier space this gives us

\[ \tilde{a}^\mu(k) = \eta \xi \left( x^\mu(k) - \frac{k(k \cdot x^\mu(k))}{|k|^2} \right) + \eta \left(1 - \xi\right) \frac{k(k \cdot x^\mu(k))}{|k|^2}, \quad (4.4) \]

where \( \tilde{a}^s(k) \) or \( \tilde{a}^c(k) \) are the Fourier sine and cosine transformed acceleration fields (i.e. \( \mu \) is \( s \) or \( c \)), \( \xi \) sets the relative ratio of the solenoidal and compressive components and

\[ \eta = \sqrt{\frac{3}{1 - 2\xi + 3\xi^2}} \quad (4.5) \]

renormalises the total amplitude. \( \xi \) can run from 0 for fully compressive forcing to 1 for fully solenoidal forcing. It should be noted that because compressive forcing can produce solenoidal modes, there is no simple scaling between the driving mix parameter, \( \xi \), and the resulting turbulent forcing parameter, \( b \) (see Chapter 1), of the gas (see e.g. Federrath et al. 2008, and references for suggested relations).

The value of the real space acceleration field at position \( r \) is then

\[ a(r) = 2f_{\text{tot}} \left( \sum_k f_k \left( \tilde{a}^c(k) \cos(k \cdot r) - \tilde{a}^s(k) \sin(k \cdot r) \right) \right), \quad (4.6) \]
where the sum is over all the modes, $f_k$ is used to adjust the relative amplitude of a given mode and $f_{\text{tot}}$ is a tuning parameter used to adjust the total strength of the forcing.

In a periodic box of side length $L$, the wavenumbers that can fit into the box go as $k_n = 2\pi n / L$. Considering only wavevectors that have positive components in the $k_x$ direction (the acceleration field is real, so we only need one half of phase-space), for a given wavevector with positive components $\mathbf{k} = (k_x, k_y, k_z)$ there also exist wavevectors in the other three remaining octants: $(k_x, -k_y, k_z), (k_x, k_y, -k_z)$ and $(k_x, -k_y, -k_z).$ Thus for every mode, there are 4 possible wavevectors. To drive over a range of spatial frequencies, we include modes such that $k_{\text{min}} \leq |\mathbf{k}| \leq k_{\text{max}}$. If there is a large difference between the size of the periodic box and the driving scale, this can result in an unfeasibly large number of modes to drive. This becomes particularly problematic if we wish to drive at ISM scales in a galaxy-scale simulation. For example, driving at 20 pc in a 100 kpc box corresponds to $n = 5000$, while driving at 30 pc corresponds to $n = 3333$ (to the nearest integer). In 1D, driving between 20–30 pc corresponds to driving 1667 modes. However, in 3D, we must drive modes that fit into the hemispherical shell in $k$-space described by $k_{\text{min}} \leq |\mathbf{k}| \leq k_{\text{max}}, \ k_x \geq 0$. This results in a total of $\sim 1.8 \times 10^{11}$ modes. Each one of these must be generated every time the acceleration field is updated. This is clearly intractable and we do not actually require such fine sampling, particularly of the longer wavelengths. We therefore calculate the acceleration field in a much smaller periodic box and tile it throughout the domain. For example, driving across this same range of spatial frequencies using a 0.5 kpc box means that only 24584 modes need to be sampled.

We can also selectively choose which gas is driven by the field. For example, to selectively drive in star forming regions, we set a threshold density equal to the star formation threshold above which gas cells will receive the acceleration from the driving field at their position. For the purposes of this demonstration, we drive between 20–30 pc in gas above our star formation threshold of 100 cm$^{-3}$. We inject at equal amplitudes across all modes (i.e. $f_k = 1$ for all $|\mathbf{k}|$). We could use $f_k$ to drive with a particular power spectrum (e.g. corresponding to Kolmogorov or Burgers turbulence) but because we are driving across such a narrow band we choose a flat spectrum and allow the turbulence to cascade down to smaller scales. We use a correlation time of $t_s = 0.7$ Myr and tune the amplitude of the turbulent driving such that in a run with an equal mix of solenoidal and compressive forcing and no SNe it increases the steady-state scale height of the disc by a factor of two relative to a run without turbulent driving. Quantitative results will of course
require a more careful calibration of the parameters to more physically motivated values.

### 4.2.3 Initial conditions

We use the same initial conditions as the fiducial galaxy from Chapter 2. We repeat the relevant details here. The total mass of the galaxy is $10^{10} M_\odot$. The dark matter halo is modelled with a static background potential corresponding to an NFW profile (Navarro et al. 1997) with a concentration of $c = 10$. We include a stellar disc and bulge in the initial conditions with masses of $1.75 \times 10^8 M_\odot$ and $1.75 \times 10^7 M_\odot$, respectively. The disc is exponential in radius with a scale length of 0.7 kpc and Gaussian in height with a scale height of 0.07 kpc. The bulge follows an Hernquist (1990) profile with scale length 0.07 kpc. The star particles present in the initial conditions do not contribute to the feedback. We initialise a gas disc with a mass of $1.75 \times 10^8 M_\odot$ with the same surface density profile as the stellar disc. The vertical profile of the gas disc is set so that it achieves hydrostatic equilibrium using the initial gas temperature of $10^4$ K. The gas and stars initially have a metallicity of $0.1 Z_\odot$. The disc and the bulge are generated using MAKENEWDISK (Springel et al. 2005). The angular momentum of these components are set such that they have the same specific angular momentum as the whole system, assuming a spin parameter of $\lambda = 0.04$. We include a uniform background CGM which fills the box with a density of $10^{-6}$ cm$^{-3}$, a temperature of $10^6$ K and no metals. We use a gas cell and star particle mass resolution of 2000 $M_\odot$ (the lowest resolution from Chapter 2).

### 4.3 Results

We perform 8 simulations: half are run without SNe (star particles still inject mass and metals corresponding to SNe, but no momentum is injected) and half are run with SNe. We perform runs without turbulent driving and with fully solenoidal ($\xi = 1$), mixed ($\xi = 0.5$) and fully compressive ($\xi = 0$) turbulent driving. Fig. 4.1 shows face-on and edge-on projections of the gas at 250 Myr. Without turbulent driving or SNe, the disc contains extremely dense clumps. It is quite likely that some degree of artificial fragmentation has been experienced. The total surface density at 250 Myr is substantially lower than the initial conditions because most of the gas has been converted into stars while the disc itself is very thin. The morphology of the gas disc with the inclusion of turbulent driving is essentially the same for all three values
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Fig. 4.1 Gas density projections for the 8 simulations, face-on and edge-on at 250 Myr. The top row shows the simulations without SN feedback while the bottom row shows the simulations with SNe. The columns show different types of driving, from left to right: no driving, purely solenoidal driving, mixed driving, purely compressive driving. Without turbulent driving or SNe, the gas fragments into dense clumps. With the addition of SNe, these clumps cause clustered feedback that disrupts the disc. Turbulent driving prevents the formation of clumps, while the addition of SNe leads to more diffuse gas and a thicker disc.

of $\xi$. The disc does not form such dense clumps as the simulation without driving, instead producing more extended clouds and filamentary structures. This suggests that the turbulent driving has provided pressure support that has prevented the disc from fragmenting to the same extent. Viewed edge on, the disc appears to be thicker (by design). Fig. 2.21 in Appendix 2.D shows a simulation using the same initial conditions and mass resolution but using a non-thermal pressure floor that ensures the Jeans length is resolved by 6.35 cell diameters. The pressure floor is also able to prevent artificial fragmentation to a similar degree, but has a slightly smoother morphology.
Without turbulent driving, dense clumps initially form in the simulation with SNe before the resulting concentrated feedback disrupts the disc. The feedback is much stronger than what was seen in Chapter 2 without the pressure floor. The degree of fragmentation results in an effectively higher star formation efficiency leading to correspondingly more violent feedback. This produces a strong outflow mass loading factors ranging between 10–100 at 1–10 kpc from the disc midplane. By 250 Myr, some gas has returned to the centre resulting in a diffuse spiral structure when the disc is viewed face-on and a complex vertical structure edge on. The three simulations with SNe and turbulent driving all have similar morphologies to each other. Again, like the equivalent simulations without SNe, a complex structure of clouds and filaments is apparent when viewed face on. The structures contain fewer dense knots of gas than the no feedback runs and there is more diffuse gas present outside of the dense structures. Viewed edge on, the addition of feedback has resulted in a slightly thicker disc and more diffuse material present above the disc plane. Much like the corresponding simulations in Chapter 2, no outflows of any substance are generated with gas circulating within a few kiloparsecs of the disc plane with mass loading factors between 0.1–0.5.

Fig. 4.2 shows the mass in newly formed stars and the SFRs as a function of time for the 8 simulations. Without turbulent driving, the SFR quickly rises to $1 M_\odot \, yr^{-1}$. Most of the available gas in the disc is converted into stars by 250 Myr, the SFR beginning to drop slightly from 100 Myr as the gas supply is exhausted and the surface density of the disc drops. With the inclusion of turbulent driving, SFR is initially suppressed. Fully solenoidal forcing suppresses the SFR by slightly under two orders of magnitude at 50 Myr compared to the simulation without driving and an order of magnitude more than the run with fully compressive turbulence. Mixed forcing lies between the compressive and solenoidal cases. From 50 Myr to 170 Myr, the simulation with fully compressive turbulence has a SFR that is increasing very slightly as the disc structure clumps up. The SFR drops from 170 Myr. This is again due to depletion of gas reservoir; the transition occurs at approximately the same amount of stellar mass formed as the no driving simulation had produced when its SFR started dropping. The simulation with solenoidal and mixed forcing have a SFR that rises steeply from 50 Myr as the gas clumps, eventually forming the same mass of stars as the compressive forcing simulation by 250 Myr.

Without turbulent driving, the simulation with SNe initially follows the same evolution as the run without SNe, experiencing a rapid increase in SFR. At 30 Myr, the resulting rate of SNe is very high. This, combined with the clustering of star formation, leads to strong feedback which suddenly suppresses the SFR. The
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Fig. 4.2 Newly formed stellar mass (top) and SFRs (bottom) as a function of time for the 8 simulations. Without feedback or turbulent driving, the SFR rapidly climbs to $1 \, M_\odot \, yr^{-1}$, only falling as it consumes the gas supply. Compressive turbulent driving initially drops the SFR by a factor of a few, while solenoidal driving suppresses it by an order of magnitude (mixed driving lies in the middle). Eventually, the rates rise and they all form the same stellar mass by 250 Myr. Without turbulent driving, the SN are highly efficient at quenching the SFR because of the increased clustering of stars and the initially high SFR. Turbulent driving allows the SN feedback to produce a steady SFR a factor of a few lower than the no feedback simulations. Again, SFR is suppressed more by solenoidal driving than compressive.

SFR begins to rise again gradually from 180 Myr as gas begins to build up again. The introduction of turbulent driving results in very steady SFRs. Unlike the simulations without SNe, these do not increase much over the course of the run. Again, solenoidal turbulence provides the greatest suppression of SFR. It actually produces a similar stellar mass at 250 Myr to the no driving case without having an extreme burst of feedback.

Fig. 4.3 shows the evolution of the disc scale height with time. We define the
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Fig. 4.3 The time evolution of the disc scale height, defined as the height enclosing 63% of the gas mass, for the 8 simulations. The disc flattens substantially over first 50 Myr as the gas cools rapidly. Without driving or feedback, it remains at a few tens of parsecs. Driving approximately doubles the scale height. Without driving, SNe efficiently disrupt the disc and the resulting outflow manifests as a rapid increase in scale height. The addition of SN to driving results in a factor of a few increase in scale height relative to the no feedback simulations.

The scale height as the height above (and below) the disc that encloses 63% of the mass\(^1\) i.e. the mass within \(h\) if the gas had an exponential profile. This is somewhat arbitrary, but the results are insensitive to reasonable variations of the enclosed mass fraction. All the simulations experience a rapid vertical compression over the first 30 Myr as the gas cools and loses pressure support. Without turbulent driving, the scale height drops to approximately 10 pc (further contraction is essentially halted by the lack of force resolution). It rises slowly as the simulation progresses. This is an effect of the rapid gas consumption moving the relative scale height upwards rather than a physical thickening of the disc. Runs with driving follow the same pattern, but the collapse is held in a steady state at roughly 20 pc (the driving was tuned to achieve this). The different types of forcing produce roughly the same result. With the inclusion of SNe (but no driving), the scale height rapidly increases after the initial collapse as a result of the strong outflow. In the simulations with driving and SNe, the scale height is a factor of a few above the no feedback driving simulations. The slight differences apparent between the different types of driving are a result of the different SN rates caused by the different SFRs

\(^1\)We actually define the total mass as that enclosed within 10 kpc. The enclosed mass profile has essentially converged by this height because the CGM makes up a negligible mass.
Fig. 4.4 Density PDFs of the gas within a radius of 6 kpc and a height of 1 kpc from the disc plane for the 8 simulations at 50 Myr (top) and 250 Myr (bottom). Without driving or feedback, a high density tail is apparent. The addition of feedback removes this feature and eventually pushes the peak of the distribution to lower densities. Turbulent driving also removes this tail initially, but without feedback it reemerges to some degree. The inclusion of feedback in combination with turbulent driving is more efficient at suppressing the high density tail.

We show density PDFs of the gas density at 50 Myr and 250 Myr for the 8 simulations in Fig. 4.4. As is apparent from Fig. 4.2, the simulation without turbulent driving or feedback produces a high density tail. Adding SNe has substantially reduced this tail by 50 Myr and substantially suppressed gas above a few $1 \text{ cm}^{-3}$ by 250 Myr. At 50 Myr, the runs with turbulent driving show a peak at $100 \text{ cm}^{-3}$, whereupon the PDF rapidly drops. This is a form of artificial self-regulation because the turbulent driving is applied from this density upwards. Above this density, the turbulent driving runs suppress dense gas to the same extent or better than the SN simulation without driving. The compressive driving simulation has
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a larger high density tail than the runs with solenoidal forcing. While a broader density distribution is in qualitative agreement the predictions for turbulent forcing described in Chapter 1 for isothermal turbulence (see eq. 1.46 and Fig. 1.6), this is more likely a result of the solenoidal driving being more effective at stabilising dense material against collapse rather than the compressive driving forcing gas to higher densities. By 250 Myr, a dense tail has built up (although it is less prominent in the simulation with compressive driving because of gas consumption). If gas becomes too dense and structures form on scales substantially below the injection scale of the turbulence, they are likely to become impervious to turbulent disruption. If the driving is unable to completely stabilise the gas from global collapse, this will lead to increased clumping over time. This issue in particular will be related to the scale on which we dissipate turbulence and will be resolution dependent. The gradually increasing metallicity of the gas throughout the run and the increased cooling rates may also play a role. The simulations with SNe and turbulent driving behave similarly to the simulations without SNe, albeit with greater suppression of dense gas.

4.4 Conclusions

We have presented a new scheme for the idealised modelling of ISM-scale turbulence in galaxy-scale simulations and provided a basic demonstration of the technique. We find that we are able to stabilise gas against extreme fragmentation without smoothing the whole disc morphology. However, without SNe high density gas is gradually created over 250 Myr, leading to gradually rising SFRs. Without injecting turbulence (or using a pressure floor), the resulting fragmentation and highly clustered star formation gives rise to a violent burst of SN feedback which disrupts the gas disc. The inclusion of turbulent driving leads a steady state, with well regulated SFRs. With the exception of the explosive feedback experienced in the simulation with SNe but no turbulent driving, we are unable to launch significant outflows. This is consistent with our findings in Chapter 2 about the resolution requirements for outflow generation. Finally, we find that driving solenoidal turbulence results in a greater suppression of SFRs than compressive turbulence.

While this technique shows promise for future studies, further work has to be carried out to allow physically meaningful comparisons with observations to be made. The results are likely to be very resolution dependent, particularly as turbulence is currently numerically dissipated at the resolution limit. We also need to calibrate our turbulent driving amplitude to match observed levels of turbulence
in GMCs. While it is possible to produce analytic estimates for the choice of parameters that should be adopted in the driving model to produce a given level of turbulence, empirically we find that a substantial level of tuning is required. In order to facilitate this tuning and provide a method to compare our ISM kinematics with observations, we need to characterise the turbulence on the GMC scale. While it is relatively simple to measure relevant quantities such as velocity dispersions in a global sense, comparing to observed GMC properties and scaling relations (e.g. Larson’s laws) will require more sophisticated techniques such as the use of some form of clump-finding algorithm.
5. Conclusions

In this thesis, high resolution hydrodynamic simulations were used to study the complex physics that drive and regulate star formation in galaxies, with a particular emphasis on modelling supernova feedback and its effects on galaxy properties. Remarkable progress has been made over the last few decades in our understanding of physical cosmology and cosmic structure formation, driven by advances in instrumentation, observational techniques, analytic theory and numerical modelling. Having said this, there are still many aspects of galaxy formation that remain uncertain. In particular, the radiative cooling and accretion of gas into the centre of dark matter haloes, the creation of a multiphase interstellar medium, the subsequent formation of stars and the resulting back-reaction on the gas from stars in the form of radiation, winds and SNe constitute an interlinked set of mechanisms that determine a wide range of galaxy and extragalactic properties. These include the stellar masses, morphologies, structures, kinematics and chemistry of galaxies and their host haloes, as well as the ability of galaxies to ionize and metal enrich the circumgalactic medium.

The presence of a multitude of baryonic processes interacting in an intricate fashion over many orders of magnitude in scale presents a challenge to the development of theoretical models. Numerical simulations provide a key tool to unpicking the details of how galaxies are formed and evolve. However, even with modern advances in computing, simulating the full dynamic range, from large-scale structure to the sites of star formation itself, is typically beyond reach. This leads to the adoption of subgrid models, designed to approximate unresolved physics. The level of abstraction adopted by the scheme is dependent on the scale of the simulation in which it is deployed. For example, large scale cosmological simulations typically have a resolution several orders of magnitude larger than the scales on which star formation and stellar feedback operate. Thus, subgrid models must take a phenomenological approach, making use of tunable parameters to replicate observations. This limits their predictive power on small scales. Alternatively, simulations of individual galaxies, either in isolated or cosmological zoom-in simulations, allow a much higher resolution to be adopted, such that the processes of star formation and stellar feedback can be modelled in a more physically motivated
In Chapter 1, I began by briefly reviewing relevant concepts from cosmology and cosmic structure formation. I introduced the various numerical methods employed in contemporary galaxy formation simulations. I then summarised the current observational and theoretical understanding of star formation in galaxies, paying particular attention to the way in which star formation efficiencies are regulated on different scales, leaving a detailed discussion of stellar feedback processes for subsequent chapters.

In Chapter 2, I presented my new implementation of star formation and SN feedback in the moving-mesh code Arepo. This is specifically designed for simulations that achieve a high enough resolution to begin treating these processes in a relatively explicit manner, with the aim of developing physically motivated, ab initio, models. The scheme has several key properties. Firstly, it individually time-resolves SNe, stochastically sampling from tabulated SN rates that are dependent on time and metallicity. The spacing between star formation and subsequent SN events can have a strong impact on the resulting efficiency of the feedback. The properties of star forming clouds can significantly change (for example, the density can increase) over the intervening 3–40 Myr. Detonating SNe immediately after the star particle is formed can also artificially curtail star formation by preventing other nearby gas from forming stars when it would otherwise have done so in the next few Myr. In addition to the offset in time between star formation and SN impacting subsequent evolution, the offset between individual SNe in the same region is potentially more important. Some subgrid SN schemes inject the total feedback budget for their star particles simultaneously (sometimes tens or hundreds at once). While this helps overcome numerical overcooling caused by poor resolution, it artificially increases the efficiency of feedback. Making sure that SNe are realistically spaced in time ensures that any increase in the efficiency of the feedback as a result of multiple SNe contributing to ‘superbubbles’ is physically correct. My scheme allows subsequent SNe to benefit from the clearing of gas by previous SN in a realistic manner and the degree to which this occurs is to some extent a quantity that is predicted by the simulations, rather than being arbitrarily enforced by the subgrid scheme.

Secondly, I have adopted a weighting scheme that ensures an isotropic distribution of feedback quantities into surrounding gas. While generally not a problem for grid codes (with adequate resolution), the spatial distribution of resolution elements in a Lagrangian code is not typically homogeneous by design. A major advantage of such a method is that it has a smoothly varying spatial resolution as a function
of density. The downside is that when distributing feedback quantities to nearest neighbours around a star particle, denser regions will receive preferentially more. In the presence of strong density gradients this can lead to significantly anisotropic behaviour, particularly in combination with a momentum based feedback scheme. In Appendix 2.B I demonstrated that a traditional SPH kernel and mass weighting scheme leads to unphysical coupling of feedback momentum into the plane of a galactic disc. My feedback scheme guarantees an isotropic distribution of feedback quantities (in all but the most extreme cell distortions) by making explicit use of the local Voronoi mesh geometry. This represents a significant departure from previous techniques adopted in AREPO.

The SN feedback scheme is capable of depositing mass, metals, energy and/or momentum into local gas in a variety of different ways, covering some of the typical methods used in the literature. These range from simple injections of thermal or kinetic energy (or a combination), the use of delayed cooling to enforce a subsequent adiabatic phase on the gas or a mechanical scheme that injects the correct amount of momentum corresponding to the phase of the SN remnant resolved (with a density and metallicity dependence). The scheme was deployed as part of an in depth study into the effects of different subgrid SN schemes on the results of galaxy formation simulations. This took the form of isolated simulations, the bulk of which were of a system of virial mass $10^{10} M_\odot$. The simulations were carried out at three resolutions spanning two orders of magnitude in cell mass to test convergence properties, the highest of which was capable of resolving the Sedov-Taylor phase of SN blastwaves. As expected, the simple injection of thermal or kinetic energy lead to overcooling at all but the highest resolution. The delayed cooling schemes proved very difficult to calibrate a priori, resulting in unphysically strong feedback. The mechanical scheme was able to provide convergent SFRs over all three resolutions, was in good agreement with the Kennicutt-Schmidt Law and preserved the disc structure. It also agreed with the simple energy injections at the highest resolution, suggesting it converges to the correct answer. While at the highest resolution, the scheme produced strong, multiphase outflows with mass loading factors between 1–10, it was unable to produce realistic outflows at lower resolutions. This may be caused by an oversimplification of the modelling of the SN remnant evolution; for example, this assumes a uniform background medium. Alternatively, the efficiency of the SNe may be enhanced by other forms of stellar feedback, such as photoionization and radiation pressure. It is also possible that a more sophisticated scheme to model star formation could lead to more clustered SN and stronger feedback.

While isolated galaxy simulations provide controlled conditions to understand
the details of the various physics at play in a clean manner, their predictive power is limited by the lack of a realistic cosmological environment. In Chapter 3, I applied my mechanical SN feedback scheme to high resolution cosmological zoom-in simulations of five $z = 4$ dwarfs with virial masses between $2 - 6 \times 10^9 \, M_\odot$. It is often assumed that in this mass range, the supply of feedback energy provided by SN alone is able to suppress SFRs sufficiently to match the extremely low stellar masses predicted by abundance matching. However, SN feedback proved to be extremely inefficient in all but one galaxy. This is caused by the build up of dense gas, often associated with a merger, that overwhelms the ability of the SNe to clear material. This is a physical effect as opposed to classical numerical overcooling. The scheme continues to inject the correct amount of momentum for a SN occurring in gas of that density. This suggests that other mechanisms (such as other sources of stellar feedback) may be required to prevent the catastrophic collapse of gas at the centre of haloes prior to the first SNe occurring. At these halo masses, the UV background is unable to prevent the build up of dense gas (with the caveat that it is modelled as time-varying but spatially uniform). I demonstrated that these results are robust to reasonable changes of the star formation prescription, increases of resolution and modifications to the UV background and pressure floor. Only by taking an extreme choice of star formation prescription parameters (setting the efficiency per free-fall time to 100%) were the resulting galaxies near the correct stellar mass. The extent to which this choice is physical is questionable.

In Chapter 4, I presented a new scheme to model the effects of a turbulent ISM in isolated galaxy simulations by idealised artificial driving of turbulence. The turbulence is injected into a narrow range of spatial frequencies on the scale of star forming clouds. This compensates for missing physics that would drive ISM turbulence in nature. While obviously limiting the predictive power of the simulations in that regard, it allows the study of the impact of a turbulent ISM on the fragmentation of the galactic disc, the spatial distribution of star formation and subsequent feedback, and the influence on galaxy properties. I demonstrated the new scheme with a highly simplified setup, showing that with turbulent driving and SN it is possible to achieve a steady SFR and a clumpy morphology of filaments and clouds. I discussed the directions that future development of this technique would take.

I have demonstrated in this thesis that a careful analysis of subgrid schemes is necessary in order to disentangle real physical insight from numerical bias. Given the proliferation of complex schemes to model a variety of physical processes (from star formation prescriptions to the multitude of feedback sources), in depth studies
into their numerical properties and, in particular, the way in which they interact with each other will be key to producing robust predictions from theory that can be compared to observations. When a large number of processes are modelled simultaneously, it is possible that physically realistic results can be produced by converging on to the wrong balance of physics.

Having dealt with SN feedback in some depth, it is clear (particularly from Chapter 3) that all sources of stellar feedback likely play some role in regulating star formation and galaxy properties in general, even in low mass haloes. The study of the manner in which they interact (which will be complicated and non-linear) is one direction in which future advances in galaxy formation theory will occur. Indeed, significant progress along these lines has already been made in recent years (see e.g. Hopkins et al. 2014, 2017; Hu et al. 2017). In addition, the development of sophisticated subgrid schemes with reference to higher resolution, smaller scale simulations of ISM processes provides an interesting avenue with which to couple scales and produce physically motivated models. This is especially applicable to treatments of star formation. From a numerical perspective, more direct modelling of radiation feedback with radiation-hydrodynamic simulations is becoming increasingly tractable, providing a more accurate picture of how this affects galaxy evolution (see e.g. Rosdahl et al. 2015; Kimm et al. 2017). The same is true of another potential source of feedback, cosmic rays, which requires modelling of magnetic fields (see e.g. Pakmor et al. 2016; Simpson et al. 2016; Pfrommer et al. 2017).

Of course, the validity of theories of galaxy formation must be tested against observations. The wealth of data from current and upcoming surveys (e.g. DES, VISTA, GAIA, LSST) and deep imaging instruments (HST, JWST and in the next decade ELT, GMT, TMT) will provide more detailed constraints on galaxy properties that can guide theoretical work, particularly for local dwarfs and high redshift analogs which are particularly sensitive to the details of stellar feedback. On the small scale, ALMA has provided and will continue to provide insight into small scale star formation physics that will be necessary to improve modelling in this area.

With increasing computational power, simulations that are able to resolve the ISM to at least some extent are becoming more common (at least in low mass haloes and/or high redshift). This provides the opportunity to treat the complex baryonic physics active on those scales in a physically motivated and consistent manner. The ability to couple such a wide range of scales will provide a crucial step on the road to fully ab initio models of galaxy formation.
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