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Steady outflows in giant clumps of high-*z* disc galaxies during migration and growth by accretion

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ABSTRACT

We predict the evolution of giant clumps undergoing star-driven outflows in high-z gravitationally unstable disc galaxies. We find that the mass-loss is expected to occur through a steady wind over many tens of free-fall times ($t_{\rm ff} \sim 10\,$Myr) rather than by an explosive disruption in one or a few $t_{\rm ff}$. Our analysis is based on the finding from simulations that radiation trapping is negligible because it destabilizes the wind (Krumholz & Thompson 2012, 2013). Each photon can therefore contribute to the wind momentum only once, so the radiative force is limited to L/c. When combining radiation, protostellar and main-sequence winds, and supernovae, we estimate the total direct injection rate of momentum into the outflow to be $2.5\,L/c$. The adiabatic phase of supernovae and main-sequence winds can double this rate. The resulting outflow mass-loading factor is of order unity, and if the clumps were to deplete their gas, the time-scale would have been a few disc orbital times, to end with half the original clump mass in stars. However, the clump migration time to the disc centre is of the order of an orbital time, about 250 Myr, so the clumps are expected to complete their migration prior to depletion. Furthermore, the clumps are expected to double their mass in a disc orbital time by accretion from the disc and clump–clump mergers, so their mass actually grows in time and with decreasing radius. From the six to seven giant clumps with observed outflows, five are consistent with these predictions, and one has a much higher mass-loading factor and momentum injection rate. The latter either indicates that the estimated outflow is an overestimate (within the 1σ error), that the star formation rate has dropped since the time when the outflow was launched or that the driving mechanism is different, e.g. supernova feedback in a cavity generated by the other feedbacks.

Key words: stars: formation – ISM: jets and outflows – galaxies: formation – galaxies: ISM – galaxies: spiral.

1 INTRODUCTION

In our developing picture of violent disc instability (VDI) at high redshift, the gas-rich discs fed by cosmological streams give birth to giant baryonic clumps that are the sites for intense star formation. The clumps are expected to migrate towards the disc centre on an orbital time-scale where they coalesce into the central bulge (Noguchi 1999; Immeli et al. 2004; Bournaud, Elmegreen & Elmegreen 2007; Elmegreen, Bournaud & Elmegreen 2008; Genzel et al. 2008; Agertz, Teyssier & Moore 2009; Dekel, Sari & Ceverino 2009; Ceverino, Dekel & Bournaud 2010). This was proposed as a mechanism for the formation of galactic spheroids, in parallel with the traditional scenario of spheroid formation by mergers (Genzel et al. 2008; Dekel et al. 2009) as well as a scenario for the formation of globular clusters (Shapiro, Genzel & Förster Schreiber 2010), and for feeding the central black holes (Bournaud et al. 2011, 2012). However, stellar feedback can generate outflows from the clumps (Krumholz & Dekel 2010; Murray, Quataert & Thompson 2010). These outflows were assumed to be very intense on a time-scale of a few free-fall times and thus lead to significant mass-loss and possibly to clump disruption (Murray et al. 2010; Genel et al. 2012; Hopkins et al. 2012b). Murray et al. (2010) argued that the high-z giant clumps are likely to be disrupted by momentum-driven feedback, as are their smaller counterpart molecular clouds in the Milky Way at low redshift, but Krumholz & Dekel (2010) pointed out that this would be possible only if the efficiency of star formation per free-fall time $\epsilon_{\rm ff}$ is significantly higher than the value implied by observations of both nearby and high-z galaxies (Krumholz & Tan 2007; Krumholz, Dekel & McKee 2012), namely the value associated with the Kennicutt–Schmidt (KS) relation.

Genzel et al. (2011) reported pioneering observational evidence for outflows from giant clumps in five $z \sim 2$ galaxies. The star

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formation rate (SFR) is estimated from the Hα luminosity. The clump properties of radius and characteristic velocity are measured directly, or alternatively the gas mass is derived from the SFR assuming that the KS relation observed on sub-galactic scales at $z = 0$ (Krumholz & Tan 2007; Krumholz et al. 2012) and on galactic scales at $z \sim 2$ (Daddi et al. 2010a,b; Genzel et al. 2010; Tacconi et al. 2010, 2013) also holds within individual giant clumps at $z \sim 2$ (see preliminary results by Freundlich et al. 2013). The outflow velocity is evaluated based on both the centroid blue-shift and the width of the broad-line component of the Hα emission, and the mass outflow rate is estimated using a variety of alternative models. Based on these observations, Genzel et al. (2011) estimated that the typical clumps in their sample drive winds of mass-loading factor $\eta \sim 1$, namely a mass outflow rate that is comparable to the SFR. These winds are expected to deplete the clump gas in several hundred Myr after turning about half the clump mass into stars. Two extreme cases, both in the same galaxy, ZC406690, indicate stronger outflows, with $\eta \sim 3$–7 and depletion times of (100–200) Myr with only 12–25 per cent of the original clump mass turning into stars. In the typical clumps, the actual driving force of the outflows is (2–4) L/c, where L/c is the contribution of a single-scatter radiation force, while in the most extreme outflow it is estimated to be as large as 34 L/c (though with a very large uncertainty).

Newman et al. (2012) performed follow-up observations at higher resolution on the galaxy with the extreme clumps, ZC406690. They compared two clumps in this galaxy, both of which are driving winds, but with very different properties. One of the clumps shows a considerably larger mass, energy and momentum flux than the other. They propose that these two clumps represent different evolutionary stages of the same phenomenon, and that the more energetic of the two outflows cannot easily be explained by any of the wind launching mechanisms that have been proposed in the literature.

In this paper, we seek to provide a unified framework for comparing different potential outflow launching mechanisms, and then use this framework to predict the outflow properties expected from stellar feedback, and understand what can be learned from the observations conducted to date. We consider the momentum injected into the wind by momentum-conserving (PC) stellar feedback mechanisms, and by the more energy-conserving (EC) supernova feedback. We predict the expected mass-loading factor and momentum injection efficiency. We focus in particular on the question of whether the migrating clumps arrive at the centre massive and intact or lose most of their mass to outflows while still in the disc.

In Section 2, we develop a simple theoretical framework for dealing with the momentum that drives outflows from star-forming clumps. We present in comparison the time-scale for clump migration and the accretion rate into the clumps during migration. In Section 3, we go through the momentum budget for the outflows. In Section 4, we address the implications for the evolution of high-z giant clumps. In Section 5, we compare the observational estimates to the predictions, and discuss the implications on the outflow driving mechanisms. In Section 6, we conclude our results and discuss them.

2 THEORETICAL FRAMEWORK

2.1 Momentum versus energy feedback

Newly formed stars generate outflows by injecting momentum and energy into the interstellar gas. Our goal in this section is to develop a basic theoretical machinery to describe this phenomenon. We first address the roles of momentum and energy in this context, and clarify the terminology of momentum-driven versus energy-driven feedback.

The central conceptual challenge is that cool interstellar gas is highly dissipative, so energy is always lost to radiative processes. Indeed, in cold gas the cooling time is almost always short compared to dynamical time-scales. Thus in launching a wind, what we are in the end always concerned with is the amount of momentum that is transferred by stellar feedback to the gas.[1] We can address two extreme ways for this transfer to occur. First, PC transfer, where ejecta from stars (photons, winds, supernova ejecta) collide with the interstellar medium (ISM) inelastically, transferring its momentum but losing some of its energy. Secondly, EC transfer, where stars heat the interstellar material, either radiatively or via shocks, to temperatures high enough that the cooling time becomes much longer than the dynamical time. When this happens the hot gas expands adiabatically, transferring momentum to the cool phases of the ISM as it does so. We generally refer to the former mechanism for launching a wind as momentum driven and the latter as energy driven, but this is a somewhat misleading nomenclature, because the rapid cooling in the cold phases of the ISM implies that in either case what ultimately matters is the momentum transferred to the cold gas.

The EC case is in general much more efficient. To see this, consider a source of ejecta with outflow velocity $V_{\rm s}$ ($V_{\rm s} = c$ for radiation) and mass flow rate $\dot{M}_{\rm s}$ (or the equivalent energy outflow rate for radiation). In the PC case, after time t, the ejecta has pushed a wind of mass $M_{\rm p}$ and velocity $V_{\rm p}$ obeying

$$M_{\rm p} V_{\rm p} = \dot{M}_{\rm s} t V_{\rm s}\,. \tag{1}$$

In the EC case, the ejecta has pushed a wind of mass $M_{\rm e}$ and velocity $V_{\rm e}$ obeying

$$M_{\rm e} V_{\rm e}^2 = \dot{M}_{\rm s} t V_{\rm s}^2\,. \tag{2}$$

The ratio of wind energies between the EC and PC cases is

$$\frac{E_{\rm e}}{E_{\rm p}} = \frac{M_{\rm e} V_{\rm e}^2}{M_{\rm p} V_{\rm p}^2} = \frac{V_{\rm s}}{V_{\rm p}} = \frac{M_{\rm p}}{\dot{M}_{\rm s} t} \gg 1\,, \tag{3}$$

and, more importantly, the corresponding ratio of momenta is

$$\frac{P_{\rm e}}{P_{\rm p}} = \frac{M_{\rm e} V_{\rm e}}{M_{\rm p} V_{\rm p}} = \frac{V_{\rm s}}{V_{\rm e}} = \left(\frac{M_{\rm e}}{\dot{M}_{\rm s} t}\right)^{1/2} \gg 1\,. \tag{4}$$

Both ratios are much larger than unity as long as the wind mass is much larger than the mass of the direct ejecta from the source, and the wind velocity is much smaller than the original ejecta velocity. From equation (4), we learn that the efficiency of injecting momentum into the wind, $\psi = MV/\dot{M}_{\rm s} t V_{\rm s}$, is much larger in the EC case than in the PC case. It turns out that radiative-pressure and stellar winds are PC mechanisms, while the expanding supernova ejecta have an early adiabatic phase that makes them closer to EC.

It is worth pausing for a moment to elaborate on how it is possible for the radial momentum imparted to the wind to greatly exceed that provided by the source when the flow is EC, and what distinguishes the EC and PC cases. The characteristic signature of the EC case is the presence of some mechanism that carries information and forces between different parts of the expanding shell of swept-up gas, allowing them to push off one another and thereby greatly increase their radial momenta while leaving the vector momentum

[1] Note that when considering a spherical outflowing shell one refers to the overall momentum in the radial direction, which is not necessarily conserved.

of the shell as zero. The mechanism responsible may be sound waves travelling through hot gas that communicate forces via gas pressure, it may be photons bouncing from one side of a spherical shell to the other that carry information via radiation pressure or it may be something else, as long as that something allows forces to be transmitted from one side of the expanding shell to the other. In contrast, the distinguishing feature of PC flows is that there is no causal communication between different parts of the expanding shell, and as a result they cannot push off each other and increase their radial momenta.

2.2 Momentum injected by star formation

Star formation in clumps, either giant clumps at $z \sim 2$ or much smaller star-forming clumps in the nearby Universe, generally does not continue until it has exhausted the available gas supply into stars. Instead, it ends after some fraction ϵ_* of the initial gas has been converted to stars,

$$\epsilon_* \equiv \frac{M_*}{M_c} . \tag{5}$$

The remaining gas is either directly expelled by the winds, or is removed by tidal stripping or simply drifts off after the winds have eroded the clump mass enough to render it unbound. The quantity ϵ_* is commonly referred to as the *star formation efficiency*. The time of final *gas depletion* is $t_{dep} = M_*/\dot{M}_*$, where $\dot{M}_*$ is the SFR and M_* is the final stellar mass at t_{dep}. This expulsion can occur via a gradual wind that removes mass from the clump continuously as it forms stars, via an explosive event that removes the bulk of the mass on a time-scale comparable to or smaller than the clump dynamical time, or some combination of the two. Our goal in this section is to develop some basic theoretical machinery to describe this phenomenon.

Let M_c and R_c be the mass and radius of a star-forming clump, and let

$$t_{ff} = \sqrt{\frac{{R_c}^3}{G M_c}} = \frac{G M_c}{{V_c}^3} , \quad V_c = \sqrt{\frac{G M_c}{R_c}} \tag{6}$$

be the corresponding free-fall time and characteristic velocity, respectively, where we are dropping factors of order unity for simplicity.[2] The crossing time R_c/V_c is of the same order as t_{ff}, and the escape speed is of the same order as V_c.[3] The instantaneous SFR is written as

$$\dot{M}_* = \epsilon_{ff} \frac{M_c}{t_{ff}} = \epsilon_{ff} G^{-1} {V_c}^3 . \tag{7}$$

The corresponding SFR time-scale is

$$t_{sfr} \equiv \frac{M_c}{\dot{M}_*} \simeq {\epsilon_{ff}}^{-1} t_{ff} . \tag{8}$$

[2] The expression $t_{ff} = \sqrt{3\pi/(32\,G\rho)}$, where ρ is the mean density within the clump, is $\sqrt{8/\pi}$ times the expression in equation (6).

[3] For an object with zero pressure and magnetic field the escape speed is $\sqrt{2}V_c$, but observed molecular clouds in the local universe are only magnetically supercritical by factors of $\sim$2 (e.g. Troland & Crutcher 2008), corresponding to a reduction in the escape speed to V_c. There are no direct measurements of magnetic field strengths in high-redshift giant clumps, but numerical simulations of magnetized turbulence suggest that a turbulent dynamo can rapidly amplify an initially sub-Alfvénic field to an Alfvén Mach number of order unity (Stone, Ostriker & Gammie 1998). If this occurs in giant clumps at high z, the escape speed should be reduced similarly for them.

The quantity ϵ_{ff} is known as the *SFR efficiency* per free-fall time, or the rate efficiency, to distinguish it from the overall star formation efficiency ϵ_*. Observations of star formation in a wide variety of environments at a wide variety of redshifts strongly constrain that $\epsilon_{ff} \sim 0.01$ (Krumholz & Tan 2007; Krumholz et al. 2012). In particular, there are indications from observed CO that the same relation with the same ϵ_{ff} is also valid at $z \sim 2$ (Daddi et al. 2010b; Tacconi et al. 2013). In our toy model, we assume that the SFR is roughly constant throughout the clump lifetime. This is consistent with cosmological simulations, where the SFR in clumps does not show a systematic variation with distance from the disc centre despite the continuous clump migration inwards (Mandelker et al., in preparation). We have replaced in equation (7) the instantaneous gas mass M_g by the total initial clump mass M_c. This should be a close overestimate as long as the clump is still far from its gas depletion time. When estimating the depletion time, we will replace M_c in equation (7) by $0.5\,M_c$, which will make t_{sfr} larger by a factor of 2.

The stars that form inject momentum into the remaining gas at a rate $\dot{p}_{in}$, in the radial direction. In this section, we will not distinguish between the energy-driven and momentum-driven routes for producing this momentum. Since feedback is generally dominated by massive stars, one can appeal to the 'old stars' limit (Krumholz & Dekel 2010), where we are concerned with time-scales longer than the $\sim$4 Myr lifetime of a massive star.[4] In this limit, the number of massive stars generating feedback at a given time is simply proportional to the SFR, so we can write

$$\dot{p}_{in} \equiv V_{in} \dot{M}_* , \quad V_{in} \equiv \psi_{in} V_L . \tag{9}$$

The velocity V_L characterizes the momentum carried by the stellar radiation field,

$$\frac{L}{c} = V_L \dot{M}_* , \tag{10}$$

where L is the luminosity produced by these stars. As discussed in Section 3.1, this quantity is

$$V_L = 190\,\mathrm{km\,s^{-1}} , \tag{11}$$

corresponding to an energy production of $cV_L = 5.7 \times 10^{17}$ erg g^{-1}. The dimensionless parameter ψ_{in} is the *momentum injection factor*; it measures the multiplicative factor by which the actual injected momentum is higher than that one would obtain if stellar radiation were the only source of momentum, and if every photon were absorbed only once before escaping. For convenience, we sometimes express it in terms of an effective *trapping factor*

$$\psi_{in} \equiv 1 + f_{trap} . \tag{12}$$

2.3 Wind properties

If the momentum injected by stars is able to raise material to speeds $V_w \gtrsim V_c$, this material may be driven off the clump in a steady wind.

[4] The old stars limit almost certainly applies to giant clumps at $z \sim 2$, since these have $t_{ff} \gg 4$ Myr (Krumholz & Dekel 2010). It probably applies to giant molecular clouds (GMCs) in the local universe as well, since, although these have $t_{ff} \sim 4$ Myr, the best observational estimates of GMC lifetimes are $\sim$30 Myr (Fukui et al. 2009); this is significantly uncertain, however. The old stars limit does not apply to smaller scale structures seen in our galaxy (see Krumholz et al. 2012 for a more thorough discussion).

Observed outflows from clumps suggest that V_w/V_c is of the order of a few. Given the available supply of momentum, equation (9), the actual momentum of the wind is

$$\dot{p}_w = \psi_{ej} V_{in} \dot{M}_* , \quad (13)$$

where ψ_{ej} is the *ejection efficiency* representing the fraction of the injected stellar momentum that goes into the wind. The ejection efficiency could be lower than unity if some material is raised to speeds below V_c and thus does not escape. Similarly, if the momentum injection is spatially distributed rather than point-like, there may be some cancellation of the momenta injected at different positions, again producing $\psi_{ej} < 1$. We define for simplicity the overall *momentum efficiency factor* in driving the wind,

$$\psi_w \equiv \psi_{ej}\psi_{in} = \frac{\dot{p}_w}{L/c} , \quad (14)$$

as the factor representing the ratio of actual momentum in the wind to the momentum carried by the radiation when each photon is counted once. This quantity can be addressed observationally, as opposed to ψ_{in} and ψ_{ej}.

The wind mass flow rate can be extracted from the wind momentum and velocity via

$$\dot{M}_w V_w = \dot{p}_w = \psi_w V_L \dot{M}_* . \quad (15)$$

The *mass-loading factor* of the wind is

$$\eta \equiv \frac{\dot{M}_w}{\dot{M}_*} = \psi_w \frac{V_L}{V_w} . \quad (16)$$

The corresponding time-scale for mass-loss by outflow is then

$$t_w \equiv \frac{M_c}{\dot{M}_w} \simeq \eta^{-1} {\epsilon_{ff}}^{-1} t_{ff} . \quad (17)$$

For η of order unity, this time-scale is comparable to the SFR time-scale, equation (8).

Note from $\dot{p}_w = \dot{M}_w V_w$ that the mass-loss rate for a given momentum budget is maximized if the wind *velocity factor*

$$\nu \equiv \frac{V_w}{V_c} \quad (18)$$

is as close to unity (from above) as possible, i.e. if the ejected gas is raised to the lowest possible velocity consistent with escape. If ν is a constant determined by the stellar momentum-driven ejection mechanism and is independent of the clump escape velocity (as indicated observationally, $\nu \sim 3$, see Section 5), then η is inversely proportional to V_c, namely the wind mass-loading factor is larger for less massive clumps.

The extreme limit of this phenomenon, as considered by Fall, Krumholz & Matzner (2010) and Krumholz & Dekel (2010), is an *explosive ejection*, which occurs when the momentum or energy injection is sufficient to raise the entire gas mass to speeds $\gtrsim V_c$ in a time of the order of t_{ff}. In this case, the feedback is likely to sweep up all the material in the clump and eject it explosively, halting any further star formation. The condition for this to occur is that $\dot{p}_w t_{ff} \sim M_g V_c$, which, using equation (15), reduces to the condition that

$$\frac{\dot{M}_*}{(M_g/t_{ff})} = \epsilon_{ff} \gtrsim \frac{1}{\psi_w} \frac{V_c}{V_L} . \quad (19)$$

As pointed out by Krumholz & Dekel (2010), this condition is extremely difficult to satisfy in giant clumps. Observations constrain $\epsilon_{ff} \sim 0.01$, and, as we will see below, $V_c/V_L \sim 0.5$ for giant clumps. Thus, achieving explosive ejection requires either that $\psi_w \sim 10$–100, that ϵ_{ff} in giant clumps exceed the observationally inferred values of $\sim$0.01 by a factor of $\sim$10–100 or some combination of both.

2.4 Clump depletion

Given our calculated mass-loss rates, we can also consider the implications for the star formation efficiency and lifetime of star-forming clumps. First, consider a clump that never experiences explosive ejection, and is simply eroded by a combination of star formation and a steady wind at constant rates. The star formation efficiency at the time of gas depletion will be

$$\epsilon_* = \frac{\dot{M}_*}{\dot{M}_* + \dot{M}_w} = \frac{1}{1+\eta} . \quad (20)$$

Conversely, a clump that has no steady wind ($\psi_{ej} = 0$) and also does not satisfy the criterion for explosive disruption, equation (19), will eventually turn itself completely into stars, or will undergo explosive disruption at late times when the amount of gas is reduced to the point where the clump is no longer in the old stars limit. By this point, however, it will have already turned the great majority of its mass into stars. A more realistic scenario is that a clump experiences a steady wind during its life and its star formation efficiency is given by equation (20).

One can compute clump depletion lifetimes in an analogous manner. In the case of depletion by a steady wind, the lifetime is

$$t_{dep} = \frac{M_*}{\dot{M}_*} = \frac{M_c}{\dot{M}_* + \dot{M}_w} \simeq \frac{2}{\epsilon_{ff}(1+\eta)} t_{ff} , \quad (21)$$

where $M_* = \epsilon_* M_c$ is the final stellar mass, with ϵ_* from equation (20). For the SFR that enters the last equality of equation (21), we have replaced M_c in equation (7) by $0.5M_c$, to represent the average between the initial gas mass of M_c and the final zero gas mass at depletion, and thus refer to the characteristic SFR during the period from the onset of the wind to depletion. With $\eta \sim 1$, this is much larger than t_{ff}. In the case of explosive disruption, the depletion time of the explosive phase is simply of the order of t_{ff}.

2.5 Clump migration versus depletion

In the case of giant clumps in high-redshift galaxies, the lifetime may also be limited because after some period a clump will *migrate* into the galactic centre following angular momentum and energy loss by torques from the perturbed disc, clump–clump interaction and dynamical fraction. The time required for this to happen is (Dekel et al. 2009; Ceverino et al. 2010)

$$t_{mig} \simeq 2.1 Q^2 \delta^{-2} t_d \simeq 8\, t_d, \quad (22)$$

where

$$t_d = \frac{R_d}{V_d} \quad (23)$$

is the disc crossing time, R_d is the characteristic disc radius and V_d is the characteristic disc circular velocity. The quantity δ is the mass fraction in cold disc within the disc radius, which at the cosmological steady state is $\delta \simeq 0.33$. The Toomre parameter is $Q \sim 0.68$ for a thick disc (Goldreich & Lynden-Bell 1965; Dekel et al. 2009). The migration time is thus comparable to the orbital time at the outer disc. If we approximate $t_d \simeq 3t_{ff}$, assuming that the clumps are overdensities of $\sim$10 with respect to the background disc (e.g. Ceverino et al. 2012), we get

$$t_{mig} \simeq 24 t_{ff} \simeq 12\, \epsilon_{ff}(1+\eta)\, t_{dep} . \quad (24)$$

The true clump lifetime will be the lesser of $t_{\rm dep}$ and $t_{\rm mig}$, and the corresponding star formation efficiency will be the lesser of ϵ_* and

$$\epsilon_{*,\rm mig} \simeq \frac{\dot{M}_*}{M_{\rm c}} t_{\rm mig} = \epsilon_{\rm ff}\frac{t_{\rm mig}}{t_{\rm ff}} \simeq 24\,\epsilon_{\rm ff}\,, \qquad (25)$$

where the last equality assumes again $t_{\rm d} \simeq 3t_{\rm ff}$. This expression for $\epsilon_{*,\rm mig}$ is valid when $t_{\rm mig}$ is significantly smaller than $t_{\rm dep}$, so the approximation $M_{\rm g} \sim M_{\rm c}$ in the SFR is good. If $t_{\rm mig}$ and $t_{\rm dep}$ are comparable, then a better approximation for $\epsilon_{*,\rm mig}$ should be smaller by a factor of ~2. From equation (24), if $\epsilon_{\rm ff} \simeq 0.01$, we learn that $t_{\rm mig}$ is expected to be smaller than $t_{\rm dep}$ as long as $(1 + \eta) < 8.3$. When $t_{\rm mig}$ is the shorter time-scale, namely when η is of order unity, the clump-bound mass fraction remaining after outflow mass-loss at the end of the migration is

$$\frac{M_{\rm c,mig}}{M_{\rm c}} \simeq 1 - \frac{\dot{M}_{\rm w} t_{\rm mig}}{M_{\rm c}} \simeq 1 - \eta\,\epsilon_{*,\rm mig}\,. \qquad (26)$$

For $\eta \sim 1$ this is a significant fraction of the original clump mass.

2.6 Clump mass growth during migration

Accretion on to the clumps, including clump mergers and possibly tidal stripping of the clumps, is significant during the clump migration inwards, and may actually be the dominant effect in the evolution of clump mass. As the clump spirals in towards the disc centre it accretes matter from the underlying disc. An estimate of the accretion rate is provided by the entry rate into the tidal (Hill) sphere of the clump in the galaxy, $R_{\rm T}$,

$$\dot{M}_{\rm ac} \simeq \alpha\,\rho_{\rm d}\,(\pi R_{\rm T}{}^2)\,\sigma_{\rm d}\,. \qquad (27)$$

Here, $\rho_{\rm d}$ is the density in the cold disc (gas or young stars), $\pi R_{\rm T}{}^2$ is the cross-section for entry into the tidal sphere and $\sigma_{\rm d}$ is the velocity dispersion in the disc representing here the relative velocity of the clump with respect of the rest of the disc. The parameter α is expected to be of order unity and smaller.

The tidal or Hill radius $R_{\rm T}$ about the clump is where the self-gravity force by the clump balances the tidal force exerted by the total mass distribution in the galaxy along the galactic radial direction. If the disc is in marginal Toomre instability with $Q \sim 1$, this is the same as the Toomre radius of the protoclump patch that contracts to form the clump (Dekel et al. 2009),

$$R_{\rm T} \simeq 0.5\,\delta\,R_{\rm d}\,, \qquad (28)$$

where the clump mass is given by

$$\frac{M_{\rm c}}{M_{\rm d}} \simeq \left(\frac{R_{\rm T}}{R_{\rm d}}\right)^2\,, \qquad (29)$$

with $M_{\rm d}$ referring to the mass of the cold disc. Also when $Q \sim 1$, the disc half-height h is comparable to $R_{\rm T}$,

$$\frac{h}{R_{\rm d}} \simeq \frac{\sigma_{\rm d}}{V_{\rm d}}\,, \qquad (30)$$

and

$$\delta \simeq \sqrt{2}\frac{\sigma_{\rm d}}{V_{\rm d}}\,. \qquad (31)$$

We can now evaluate the time-scale for clump growth by accretion, $t_{\rm ac}$, using equation (27). We insert $R_{\rm T}$ from equation (28), write $\rho_{\rm d} = M_{\rm d}/(2\pi R_{\rm d}{}^2\,h)$ and use equation (30) for h to obtain

$$t_{\rm ac} \equiv \frac{M_{\rm c}}{\dot{M}_{\rm ac}} \simeq \frac{2}{\alpha} t_{\rm d}\,. \qquad (32)$$

With $\alpha \sim 1/3$ (see below), the time-scale for doubling the clump mass by accretion is roughly an orbital time, comparable to the migration time.

The parameter α represents the fraction of the mass entering the tidal radius that is actually bound to the clump. If the clump collapses from an initial patch of radius $R_{\rm T}$, the particles enter the tidal radius with a velocity distribution that is similar to that of the overall disc, namely with a standard deviation $\sigma_{\rm d}$, and a distribution of kinetic energies per unit mass about $0.5\,\sigma_{\rm d}^2$ in the clump rest frame. Using equations (28), (29) and (31), the binding potential of the clump at $R_{\rm T}$ can be crudely estimated by

$$\frac{GM_{\rm c}}{R_{\rm T}} \sim \sigma_{\rm d}^2\,, \qquad (33)$$

so a significant fraction of the particles entering the tidal radius are expected to be bound.

One can estimate α by referring to the particles that actually hit the clump, of radius $R_{\rm c} < R_{\rm T}$, and are bound there. This requirement puts an upper limit on the particle impact parameter b prior to entering $R_{\rm T}$ such that the focusing of the orbit would bring the particle into $R_{\rm c}$ with a velocity smaller than the escape velocity from the clump $V_{\rm c}$ at $R_{\rm c}$. Angular-momentum conservation yields $b \simeq R_{\rm c}V_{\rm c}/\sigma_{\rm d}$. Using equation (6) for $V_{\rm c}$ and equation (33) for $\sigma_{\rm d}$, we obtain $V_{\rm c}/\sigma_{\rm d} \simeq (R_{\rm T}/R_{\rm c})^{1/2}$. Therefore, in equation (27)

$$\alpha \simeq b^2/R_{\rm T}{}^2 \simeq R_{\rm c}/R_{\rm T}\,. \qquad (34)$$

With a typical contraction factor of $R_{\rm T}/R_{\rm c} \simeq 3$ (Ceverino et al. 2012), the estimate is $\alpha \simeq 1/3$.

For a uniform disc, the relevant density $\rho_{\rm d}$ is the mean density in the disc, $\bar{\rho}_{\rm d}$, within a cylinder of radius $R_{\rm d}$ and height $2h$, and α is the same throughout the disc. However, for an exponential disc with an exponential radius $R_{\rm exp}$, the local density $\rho_{\rm d}$ is $\simeq (0.28, 0.46, 0.70)\,\bar{\rho}_{\rm d}$ at $(3, 2, 1)\,R_{\rm exp}$, respectively. Therefore, if one uses the mean density in equation (27), the effective value of α in the outer disc could be slightly smaller than estimated above.

The clump growth rate is further enhanced by mergers of clumps as they spiral in. If the clumps contain 20 per cent of the disc mass (Dekel et al. 2009), and if we require binding when the clump centres are at a distance of $2R_{\rm c}$ from each other, an analogous estimate to equation (27) gives that the time-scale for growth by mergers is roughly $t_{\rm mer} \simeq (5/2)t_{\rm ac}$. On the other hand, tidal stripping may become more pronounced at small radii, which may slow down the mass growth rate at the late stages of the migration. We expect the time-scale for mass growth to be comparable to the crude estimate in equation (32).

The above estimates for $t_{\rm ac}$ and $t_{\rm mer}$ are indeed consistent with the findings from hydrocosmological simulations, in which outflows by stellar feedback are weak by construction. In these simulations, the clump mass is found to be roughly inversely proportional to distance from the disc centre (Mandelker et al., in preparation). When following individual clumps as they accrete, strip and merge during migration, they indeed grow in mass on a time-scale that is comparable to the migration time-scale. An effective value of $\alpha \sim 0.33$ in equation (32) seems to provide a good fit to the overall clump mass growth rate in these simulations (with negligible outflows).

3 MOMENTUM BUDGET

Having developed a basic framework for how the properties of star-forming clumps – their star formation efficiencies, lifetimes and outflows – depend on the momentum injected by stellar feedback, we now turn to various possible feedback mechanisms. Our goal is to understand the momentum injection efficiency $\psi_{\rm in}$ for each mechanism. We note that a similar budgeting exercise has been carried out by Matzner (2002) for Galactic GMCs, and our general approach will follow this, applied to a quite different context.

Unless otherwise stated, all the values below are derived using a STARBURST99 (Leitherer et al. 1999; Vázquez & Leitherer 2005) calculation for continuous star formation, with all parameters set to their default values except that we use an IMF upper limit of 120 $M_\odot$ instead of the default of 100 $M_\odot$, though the difference in most quantities is small (<10 per cent). We use 120 $M_\odot$ because it is the largest mass for which evolutionary tracks are available in STARBURST99, and because observations indicate that the IMF extends to at least this mass if not significantly higher (Crowther et al. 2010). We evaluate all quantities at a time of 100 Myr after the start of star formation, but since luminosity and all other quantities are slowly varying at times >4 Myr (which is the essence of the old stars limit), different choices of age in the range 10–300 Myr make a difference of at most a few tens of per cent. Given this uncertainty, we give all results to two significant digits only.

3.1 Stellar radiation

Stellar radiation pressure has received a great deal of attention recently as a potential mechanism for disrupting giant clumps, both in analytic models (Krumholz & Dekel 2010; Murray et al. 2010) and in numerical simulations (Hopkins, Quataert & Murray 2011; Genel et al. 2012; Hopkins, Quataert & Murray 2012a; Hopkins et al. 2012b). As stated above, our STARBURST99 calculation gives $V_{\rm L} = 190\,{\rm km\,s^{-1}}$ as the momentum budget of the direct radiation field.

Recall our definition of $\psi_{\rm in} = 1 + f_{\rm trap}$ as the ratio of the momentum actually injected into the gas to that which would be imparted by direct radiation pressure alone. Thus, $\psi_{\rm in} = 1$ corresponds to a flow that receives no momentum from any source but photons that are absorbed once and then escape. If the radiation emitted by stars is trapped by the high optical depths of a dust layer, it is possible that the trapped photons will be absorbed or re-emitted multiple times before escape. If this occurs, radiation energy can build up, and the adiabatic expansion of the radiation-dominated region can result in a larger momentum transfer to the gas. Some analytic and numerical models of radiation pressure-driven feedback assume that this effect will produce a value of $f_{\rm trap} \sim \tau$, where τ is an approximate infrared optical depth (Murray et al. 2010; Genel et al. 2012; Hopkins et al. 2012a), while others assume that radiation-driven flows are strictly momentum limited, with $\psi_{\rm in}$ never exceeding a few (Krumholz & Matzner 2009; Fall et al. 2010; Krumholz & Dekel 2010), due to radiation Rayleigh–Taylor instability (Jacquet & Krumholz 2011). This instability punches holes in the gas that allow photons to leak out, and prevent the build-up of adiabatic radiation-dominated regions.

Recent radiation-hydrodynamic simulations by Krumholz & Thompson (2012, 2013) have significantly clarified the matter. They show that, in the case of a radiatively-driven wind, the trapping factor obeys $f_{\rm trap} \approx 0.5\tau_*$, where τ_* is the optical depth evaluated using the opacity at the dust photosphere, *not* the far higher opacity found deep in the dust gas where the radiation temperature is higher, as assumed e.g. by Hopkins et al. (2012a). In the old stars limit, for an object of SFR per unit area $\dot\Sigma_*$ and gas surface density $\Sigma_{\rm gas}$, this is given by (Krumholz & Thompson 2013)

$$\tau_* = 0.01\,(cV_{\rm L})_{10}^{1/2}\,\dot\Sigma_{*,0}^{1/2}\,\Sigma_{\rm gas,0}\,, \tag{35}$$

where $(cV_{\rm L})_{10} = (cV_{\rm L})/(10^{10}\,{\rm L_\odot}/({\rm M_\odot\,yr^{-1}}))$, $\dot\Sigma_{*,0} = \dot\Sigma_*/(1\,{\rm M_\odot\,yr^{-1}\,kpc^{-2}})$ and $\Sigma_{\rm gas,\,0} = \Sigma_{\rm gas}/(1{\rm g cm^{-2}})$. The values of $\dot\Sigma_*$ and $\Sigma_{\rm gas}$ to which we have scaled are typical of observed giant clumps, as discussed below, and thus we typically have $\tau_* \ll 1$, and therefore $f_{\rm trap,\,rad} \approx 0$. We conclude that the very large trapping factors of 10−50 assumed in certain simulations (e.g. Oppenheimer & Davé 2008; Genel et al. 2012; Hopkins et al. 2012a) are unrealistic.

Note that the energy in the wind is

$$\dot E_{\rm w} = \frac{1}{2}\dot M_{\rm w} V_{\rm w}{}^2 = \frac{1}{2}\psi_{\rm w}\frac{V_{\rm w}}{c}L\,. \tag{36}$$

Thus, as long as $\psi_{\rm w}$ is of the order of a few and $V_{\rm w} \ll c$, only a small friction of the photon energy is used to drive the outflow. This is indeed PC rather than EC driving of the outflow.

3.2 Photoionized gas

In Galactic GMCs, the pressure of photoionized gas is likely the dominant feedback mechanism that limits the star formation efficiency (e.g. Whitworth 1979; Williams & McKee 1997; Matzner 2002; Krumholz et al. 2006; Goldbaum et al. 2011). Photoionization raises gas to a nearly fixed temperature of roughly 10^4 K, and that temperature is maintained by radiative heating and cooling processes. As the gas expands, it transfers momentum to the surrounding medium. By integrating over the IMF and using a similarity solution to compute the evolution of expanding H II regions, Matzner (2002) estimates that this mechanism injects momentum at a rate $V_{\rm in} \simeq 260\,{\rm km\,s^{-1}}$. While this is probably the dominant feedback mechanism for GMCs in local galaxies, it is likely to be unimportant for giant clumps, for the simple reason that such clumps have characteristic speeds significantly higher than the ionized gas sound speed, $c_{\rm i} \simeq 10\,{\rm km\,s^{-1}}$ (Krumholz & Dekel 2010). As a result, ionized gas will be unable to expand and transfer momentum to the cold gas. Recent numerical simulations confirm this conjecture (Dale, Ercolano & Bonnell 2012). We may therefore disregard this mechanism for giant clumps at high redshift.

3.3 Protostellar outflows

Protostars drive collimated hydromagnetic outflows with launch speeds comparable to the escape velocity from stellar surfaces, typically $\sim$100 km s^{-1} for protostars with radii larger than those of main-sequence stars. The wind material shocks against and mixes with the surrounding dense molecular gas. Because the environment the winds encounter is very dense, and the shock velocity is not high enough to heat material past the $\sim 10^5$ K peak of the cooling curve, the post-shock gas rapidly cools via radiation, so there is no significant adiabatic expansion phase. STARBURST99 does not include protostellar winds, so we adopt an estimate of

$$V_{\rm ps} \simeq 40\,{\rm km\,s^{-1}}\,, \tag{37}$$

from Matzner (2002).

3.4 Supernovae

From our STARBURST99 calculation, supernovae occur at a rate

$$\tau_{\rm sn}^{-1} = 0.012\,(\dot M_*/{\rm M_\odot\,yr^{-1}})\,, \tag{38}$$

carry an energy[5] $V^2 \simeq 5.8 \times 10^{15}$ erg g^{-1} and carry a momentum flux that corresponds to[6]

$$V_{\rm sn} \simeq 48\,{\rm km\,s^{-1}}\,. \tag{39}$$

This represents only a lower limit on the true momentum injected by supernovae, because the extremely high post-shock temperatures produced when supernova ejecta encounter the ISM guarantee that radiative cooling is inefficient at the early stages. As a result, supernova remnants experience an EC phase when they are small, then begin to cool radiatively only once adiabatic expansion lowers their internal temperatures sufficiently. During the adiabatic Sedov–Taylor phase and the pressure-driven snowplow phase that follows it (during which the remnant interior is partially radiative), the radial momentum carried by the swept-up material increases. This process has been studied by numerous authors (e.g. Chevalier 1974; McKee & Ostriker 1977; Cioffi, McKee & Bertschinger 1988; Thornton et al. 1998), and for a uniform medium Thornton et al. find that the asymptotic momentum of a supernova remnant is

$$p_{\rm sn} = 1.7 \times 10^{43} E_{51}^{13/14}\, n_1^{-0.25}\ {\rm g\ cm\ s^{-1}}\,, \tag{40}$$

where E_{51} is the energy of the supernova in units of 10^{51} erg and n_1 is the ambient number density in units of 10 cm^{-3}. A simple estimate for the supernova momentum budget including the adiabatic phase is simply obtained from $\dot{p}_{\rm sn}$, which is given by $p_{\rm sn}$ from equation (40) multiplied by the supernova rate $\tau_{\rm sn}^{-1}$ from equation (38). This gives

$$V_{\rm sn,adiab} = 1100\, n_1^{-0.25}\ {\rm km\ s^{-1}}\,, \tag{41}$$

assuming that the energy of a single supernova is $E_{51} = 1$ (as assumed in the STARBURST99 calculation as well). Since the momentum input from a single supernova remnant is very close to linear in E_{51}, the total momentum budget is not significantly affected by the manner in which the supernovae are clustered.

However, we caution that this calculation is for a simple, one-dimensional uniform medium. As the Krumholz & Thompson (2012, 2013) results for radiation pressure show, this assumption can be deeply misleading about how effectively energy is converted into momentum in a real three-dimensional medium where instabilities can occur. It is therefore best to regard equation (41) as representing an upper limit. Determining where reality lies in between this value and the lower limit represented by equation (39) requires numerical simulations capable of following instabilities into the non-linear phase. Although such simulations have begun to appear in the literature (Creasey, Theuns & Bower 2013), the problem remains far from fully solved.

3.5 Main-sequence and post-main-sequence stellar winds

The winds of main-sequence and post-main-sequence stars carry an energy and momentum content $V^2 \simeq 1.5 \times 10^{15}$ erg g^{-1} and

$$V_{\rm ms,dir} \simeq 140\,{\rm km\,s^{-1}}, \tag{42}$$

respectively. While they therefore carry slightly less momentum than the stellar radiation field, at least some of the winds are launched at velocities large enough that the post-shock gas may have long cooling times. (This is in contrast to the much slower protostellar outflows.) As a result, it is plausible that stellar winds could experience an adiabatic phase like supernovae, and enhance their momentum transfer that way. In this case, the momentum provided by winds will have an additional term that we can write as $f_{\rm ad} V_{\rm L}$.

The main idea of the classical stellar wind bubble model of Castor, McCray & Weaver (1975) and Weaver et al. (1977) is that $f_{\rm ad} V_{\rm L} \gg V_{\rm ms,\,dir}$. On the other hand, this mechanism will not operate if wind gas is able escape from a star-forming clump without entraining significant mass, or if it undergoes rapid cooling by mixing with cooler, dense gas that brings its temperature low enough for radiative losses to become rapid. Observations of a few nearby H II regions have been able to address this question directly by using X-ray observations to probe the energy density of the shock-heated gas. Both Harper-Clark & Murray (2009), who study the Carina Nebula, and Lopez et al. (2011), who study 30 Doradus, find that the luminosity of the X-ray emitting gas implies that the pressure exerted by this gas is weaker than that exerted by photoionized gas, a result highly inconsistent with an energy-driven flow.[7]

For our fiducial estimate in this paper, we adopt Lopez et al.'s measured mean value $f_{\rm ad} = 0.3$, so that the net amount of momentum ejected by the winds from main-sequence stars is

$$V_{\rm ms} = V_{\rm ms,dir} + f_{\rm ad} V_{\rm L} \simeq 200\,{\rm km\,s^{-1}}\,. \tag{43}$$

However, we caution that none of the observed regions have conditions close to those of high-redshift giant clumps. While stellar wind gas is momentum driven and not energy driven in the local universe, a giant clump could be considerably harder for hot X-ray gas to escape. It is therefore conceivable that, under the conditions found in high-redshift giant clumps, stellar winds represent an energy-driven feedback. In this case, it is likely that stellar wind and supernova bubbles would simply add together to produce an adiabatic shell driven by the combined effects of both. The result is to increase the adiabatic energy budget by roughly 25 per cent compared to supernovae alone. Adopting this simple estimate, we find that if

[5] Note that V and V^2 here are the energy and momentum per unit mass of stars formed, not per unit mass of stars that actually end their lives as SNe.
[6] The publicly available version of STARBURST99 does not calculate the supernova momentum flux. We have modified it to do so, using the same assumptions STARBURST99 adopts in order to compute the supernova energy injection rate and mass return, i.e. all stars with an initial mass above 8 M$_\odot$ end their lives as supernovae with identical energies of 10^{51} erg. They all leave behind as remnants 1.4 M$_\odot$ neutron stars, so the mass of the ejecta is simply the final stellar mass (smaller than the initial mass due to wind losses) minus the remnant neutron star mass.

[7] Contrary to the results of Lopez et al. (2011), Pellegrini, Baldwin & Ferland (2011) argue that the X-ray emitting gas pressure in 30 Doradus is actually higher than the ionized gas pressure. Pellegrini et al.'s results differ because they assume that the X-ray emitting gas is confined to a small volume. Since the X-ray luminosity is proportional to the emission measure of the emitting gas, which is the integral of the square of the electron density along the line of sight, if one assumes that the line-of-sight length is much smaller than the transverse size of the region being observed, the density and thus pressure that one infers for a given observed luminosity rises proportionately. For our purposes, however, this distinction is irrelevant. If one assumes that the hot gas is confined to a small fraction f_X of the observed volume, the pressure P_X inferred from a given luminosity varies as $P_X \propto f_X^{-1/2}$, but the energy content of the hot gas, which varies $f_X P_X$, falls as $f_X^{1/2}$. Thus, if Pellegrini et al.'s conjecture about the geometry is correct, that implies even more strongly that feedback from stellar winds cannot be energy driven.

winds are adiabatic, then, in conjunction with supernovae, the net momentum contribution of the winds is

$$V_{\rm ms,adiab} = 275\, n_1^{-0.25}\ {\rm km\ s^{-1}}. \quad (44)$$

Note that this estimate implicitly assumes that the temperature and cooling are determined by the significantly larger energy associated with the supernovae, so that stellar winds simply pump more energy into the adiabatic bubble without significantly affecting how it cools.

3.6 Total momentum budget

Combining all the mechanisms we have enumerated (excluding photoionized gas for the reasons stated above), we see that the momentum budget for the case of a purely momentum-driven outflow is expected to be

$$V_{\rm in} = V_{\rm rad} + V_{\rm ps} + V_{\rm sn} + V_{\rm ms} \simeq 480\,{\rm km\,s^{-1}}\,. \quad (45)$$

The corresponding trapping factor is rather small,

$$\psi_{\rm in} = 1 + f_{\rm trap} \simeq 2.5\,. \quad (46)$$

By 'purely momentum-driven outflow', we refer to the case where radial momentum is conserved and neither supernovae nor stellar winds experience an EC phase during which their radial momentum is significantly boosted. This can be interpreted as a lower limit for $\psi_{\rm in}$ from stellar feedback.

If we assume that there is a significant adiabatic phase for supernovae and main-sequence winds, and that the resulting momentum injection is near the upper limit derived in the uniform medium case, then we obtain an upper limit for the net momentum budget of

$$\begin{aligned} V_{\rm in,adiab} &= V_{\rm rad} + V_{\rm ps} + V_{\rm sn,adiab} + V_{\rm ms,adiab} \\ &\simeq 230 + 1350\, n_1^{-0.25}\ {\rm km\,s^{-1}}\,, \end{aligned} \quad (47)$$

$$\psi_{\rm in} \simeq 1.2 + 7.1\, n_1^{-0.25}\,. \quad (48)$$

For massive clumps of $t_{\rm ff} \simeq 7$ Myr this is $\psi_{\rm in} \simeq 6.4$.

The actual contribution of adiabatic supernova feedback and the corresponding value of $\psi_{\rm in}$ between the above lower and upper limits is a matter of a saturated state of fully non-linear instabilities, which should be determined by appropriate numerical simulations. The two most relevant publications to date on this subject are Hopkins et al. (2012a) and Creasey et al. (2013). For the former, if we examine the runs without the sub-grid radiation model where supernova feedback dominates, the mass-loading factors are $\eta \sim$ 1–5 and the wind terminal speeds are a few hundred km s^{-1}, implying $\psi_{\rm w}$ values of a few; combined these suggest $\psi_{\rm in} \sim 5$. Similarly, Creasey et al. (2013) report a mass-loading factor η and a wind thermalization parameter η_T (which measures the fraction of supernova energy that goes into outflow; see their equation 5), both as a function of galaxy properties. With some algebra, one can show that $\psi_{\rm w} = 5.3(\eta\eta_T)^{1/2}$, and using their fitting formulae for galaxy surface densities ∼10–100 $M_\odot$ pc^{-2}, appropriate to giant clump galaxies, gives $\psi_{\rm w} \sim$ 1–3 for supernovae alone. In summary, the numerical results of both Hopkins et al. (2012a) and Creasey et al. (2013) suggest that $\psi_{\rm in}$ is likely to be roughly halfway between our upper and lower limits.

A potential way to make the supernova feedback more effective is by having the gas density in the supernova vicinity much lower than the unperturbed density within the clump of $n_1 \geq 1$. A value of $n_1 \sim 10^{-2}$ (or 10^{-3}) in equation (48) would provide a maximum value of $\psi_{\rm in} \simeq 24$ (or 41, respectively). The question is whether the other types of PC stellar feedback could generate such a low-density regime prior to the supernova explosion. We keep this mechanism outside the scope of this paper.

4 IMPLICATIONS FOR HIGH-REDSHIFT GIANT CLUMPS

4.1 Star formation and outflows

In the preceding two sections, we developed a general framework to consider the evolution of clumps as they migrate, accrete, form stars and lose mass due to star formation feedback, and we derived an estimate for the momentum budget of the feedback that drives clump winds. We now combine these results to draw conclusions about the typical evolutionary path taken by giant clumps.

The structure and dynamics of a clump are characterized by two quantities, e.g. its characteristic velocity and its free-fall time, $V_{\rm c} \equiv 100\,{\rm km\,s^{-1}}\,V_2$, and $t_{\rm ff} \equiv 10\,{\rm Myr}\,t_{\rm ff10}$, as defined in equation (6).[8]

Following the earlier discussion, the physics of outflow from giant clumps can be characterized by three dimensionless parameters, e.g. $\epsilon_{\rm ff}$, ν and $\psi_{\rm w} = \psi_{\rm ej}\psi_{\rm in}$. These are the SFR efficiency $\epsilon_{\rm ff} \equiv 0.01\epsilon_{\rm ff,-2}$, the wind velocity with respect to the clump escape velocity, $\nu \equiv V_{\rm w}/V_{\rm c} \equiv 3\nu_3$, and the wind momentum with respect to the radiation momentum, $\psi_{\rm w} = \dot{p}_{\rm w}/(L/c)$. For the values of $\epsilon_{\rm ff}$ and $\psi_{\rm in}$ we have theoretical predictions. For the values of ν and $\psi_{\rm ej}$, unity is a lower and an upper limit, respectively, but we do not have a theoretical prediction concerning how much they actually deviate from unity. Motivated by observations (see below), we assume that these deviations are by a multiplicative factor of the order of 1 or a few. We define $\psi_{\rm w} \equiv 2.5\psi_{\rm w,2.5}$. The reference value of $\psi_{\rm w} \simeq 2.5$ may refer to the case of pure momentum-driven outflow $\psi_{\rm in} \simeq 2.5$ and maximum ejection of $\psi_{\rm ej} \simeq 1$, or to a case including adiabatic supernova and stellar-wind feedback with $\psi_{\rm in} \simeq 5$ but with some losses in the ejection, $\psi_{\rm ej} \simeq 0.5$. The maximum value of $\psi_{\rm w}$, when adiabatic supernova feedback is at its maximum and the ejection is efficient, is expected to be $\psi_{\rm w} \sim 5$.

The SFR and wind mass flow rate are

$$\dot{M}_* \simeq 2.4\epsilon_{\rm ff,-2} V_2^3\ {\rm M_\odot\ yr^{-1}}\,, \quad (49)$$

$$\dot{M}_{\rm w} \simeq 3.2\epsilon_{\rm ff,-2}\psi_{\rm w,2.5}\nu_3^{-1}V_2^2\ {\rm M_\odot\ yr^{-1}}\,. \quad (50)$$

The corresponding time-scales are

$$t_{\rm sfr} = \frac{M_{\rm c}}{\dot{M}_*} \simeq 1\,{\rm Gyr}\,\epsilon_{\rm ff,-2}^{-1} t_{\rm ff,10}\,, \quad (51)$$

$$t_{\rm w} = \frac{M_{\rm c}}{\dot{M}_{\rm w}} \simeq 1\,{\rm Gyr}\,\psi_{\rm w,2.5}^{-1} V_{\rm w,400}\epsilon_{ff,-2}^{-1} t_{\rm ff,10}\,. \quad (52)$$

The mass-loading factor, equation (16), is

$$\eta \simeq \psi_{\rm w,2.5} V_{\rm w,400}^{-1} = 1.33\psi_{\rm w,2.5}\nu_3^{-1}V_2^{-1}\,, \quad (53)$$

where $V_{\rm w} \equiv 400\,{\rm km\,s^{-1}}\,V_{\rm w,400}$. With the fiducial values adopted here for momentum-driven stellar feedback from typical clumps, one expects steady winds with mass-loading factors of order unity. The maximum value, when adiabatic supernova feedback is included, is expected to be $\eta \sim$ 2–3.

[8] The relations to the clump mass and radius are $V_2 \simeq 1.15 M_{9.5}^{1/2} R_1^{-1/2}$ and $t_{10} \simeq 0.96 R_1 V_2^{-1} \simeq 0.82 R_1^{3/2} M_{9.5}^{-1/2} \simeq 1.27 M_{9.5} V_2^{-3}$ where $M_{9.5} \equiv M_{\rm c}/10^{9.5}\,{\rm M_\odot}$ and $R_1 \equiv R_{\rm c}/1\,{\rm kpc}$. Also, $n_1 \simeq 2.2 t_{10}^{-2}$. The surface density is $\Sigma \simeq 0.21 M_{9.5} R_1^{-2}\,{\rm g\,cm^{-2}}$ and $1\,{\rm g\,cm^{-2}} \simeq 4800\,{\rm M_\odot\,pc^{-2}}$.

If the clump were allowed to deplete all its gas, the final star formation efficiency would have been

$$\epsilon_* = (1+\eta)^{-1} \equiv 0.5(1+\eta)_2^{-1} \tag{54}$$

at the clump depletion time of

$$t_{\rm dep} = 1\,{\rm Gyr}\,(1+\eta)_2^{-1}\epsilon_{\rm ff,-2}^{-1} t_{\rm ff,10}\,, \tag{55}$$

where $(1+\eta)_2 \equiv (1+\eta)/2$.

The approximate values for ϵ_* and $t_{\rm dep}$ are valid when ϵ_* deviates significantly from unity, and where clump disruption is by gradual erosion rather than sudden explosive destruction. The criterion for explosive disruption, equation (19), is simply

$$\epsilon_{\rm ff,-2}\psi_{\rm w,2.5}V_2^{-1} \gtrsim 20. \tag{56}$$

This is similar to equation 9 of Krumholz & Dekel (2010), where the considerations were qualitatively similar though not exactly the same numerically. As noted there, if the clump is a typical Toomre clump with $V_{\rm c} \sim 100\,{\rm km\,s^{-1}}$, explosive disruption occurs only if either $\epsilon_{\rm ff}$ or $\psi_{\rm w}$ are significantly larger than their fiducial values, namely either the SFR is much more efficient than implied by the local Kennicutt relation, or the momentum-driven feedback is much more efficient than available in the momentum budget evaluated above. Otherwise, with the adopted fiducial values for these quantities, explosive disruption does not occur, thus validating the steady-wind approximation used.

4.2 Clump migration

In comparison, the clump migration time is

$$t_{\rm mig} \simeq 8t_{\rm d} \simeq 260\,{\rm Myr}\,R_{\rm d,7}V_{\rm d,200}^{-1}\,, \tag{57}$$

where the disc is characterized by $t_{\rm d} = R_{\rm d}/V_{\rm d}$ with $R_{\rm d,7} \equiv R_{\rm d}/7$ kpc and $V_{\rm d,200} \equiv V_{\rm d}/200\,{\rm km\,s^{-1}}$. The fiducial values of $R_{\rm d}$ and $V_{\rm d}$ are deduced from observations at $z \sim 2$ (Genzel et al. 2006, 2008), but the disc dynamical time can also be derived from the virial radius and velocity using the virial relation and the spherical collapse model and assuming a constant spin parameter λ for haloes and conservation of angular momentum during gas collapse within the dark matter halo,

$$t_{\rm d} \simeq \lambda \frac{R_{\rm v}}{V_{\rm v}} \simeq 0.07\,\lambda_{0.07}\,0.15\,t_{\rm Hubble}\,, \tag{58}$$

which at $z = 2$, where $t_{\rm Hubble} \simeq 3.25$ Gyr, gives $t_{\rm d} \sim 33\,\lambda_{0.07}$ Myr.

The relation between the depletion time and migration time is

$$t_{\rm mig}/t_{\rm dep} \simeq 0.25\,(1+\eta)_2\,\epsilon_{\rm ff,-2}\,, \tag{59}$$

where we have assumed $t_{\rm d} \simeq 3t_{\rm ff}$ for the dynamical time-scales in the disc and in the clumps. If $t_{\rm mig} \le t_{\rm dep}$, the maximum star formation efficiency possible before the clump reaches the galactic centre is

$$\epsilon_{*,\rm mig} \simeq 0.24\epsilon_{\rm ff,-2}\,. \tag{60}$$

The clump-bound mass fraction remaining at the end of the migration is

$$\frac{M_{\rm c,mig}}{M_{\rm c}} \simeq 1 - \frac{\dot{M}_{\rm w}t_{\rm mig}}{M_{\rm c}} \simeq 1 - 0.24\eta\epsilon_{\rm ff,-2}\,. \tag{61}$$

These are for $t_{\rm mig}$ significantly smaller than $t_{\rm dep}$, namely for η of order unity. The estimate for $\epsilon_{*,\rm mig}$ is an overestimate by up to a factor of the order of 2 because we assumed here $M_{\rm g} \simeq M_{\rm c}$. For the same reason, the expression for $M_{\rm c,mig}/M_{\rm c}$ is an underestimate. With the fiducial values adopted here, the clump reaches the centre while still holding on to a significant fraction of its original mass, and most likely still gas rich.

4.3 Clump mass evolution

Equation (32), with $t_{\rm d} \simeq 3t_{\rm ff}$, yields

$$t_{\rm ac} \simeq 0.18\,{\rm Gyr}\,\alpha_{0.33}^{-1}\,t_{\rm ff,10}\,, \tag{62}$$

where $\alpha_{0.33} = \alpha/0.33$. Thus, with $\alpha = 0.33$, the time-scale for doubling the clump mass by accretion is $\sim 6t_{\rm d} \sim 0.18$ Gyr, which is slightly smaller than the migration time, $t_{\rm mig} \sim 8t_{\rm d} \sim 0.24$ Gyr.

With the fiducial values for momentum-driven winds, the mass growth rate as estimated in equation (62) is faster than the outflow rate and the SFR, equations (51) and (52), implying that the accretion more than compensates for the mass-loss by outflows, making the clumps actually grow in mass as they migrate inwards. This implies in particular that the adopted estimate of migration time remains a good approximation and may even be an overestimate.

The evolution of clump mass $M(t)$ under accretion and outflows, starting from an original mass $M_{\rm c}$ at $t = 0$, is governed by

$$\dot{M} = \dot{M}_{\rm ac} - \dot{M}_{\rm w}\,. \tag{63}$$

What makes the integration of this equation simple is that the two terms on the right-hand side both scale with $M/t_{\rm ff}$. First,

$$\dot{M}_{\rm ac} \simeq \frac{\alpha}{2}\frac{t_{\rm ff}}{t_{\rm d}}\frac{M}{t_{\rm ff}}\,, \tag{64}$$

where α and $t_{\rm ff}/t_{\rm d}$ are approximated as constants, the latter being determined by the clump collapse factor from the original protoclump patch in the disc. Second,

$$\dot{M}_{\rm w} \simeq \eta\,\epsilon_{\rm ff}\,f_{\rm g}\frac{M}{t_{\rm ff}}\,, \tag{65}$$

where $f_{\rm g}$ is the star-forming gas fraction in the clump, approximated as constant. Integrating, we obtain

$$M(t) = M_{\rm c}\,{\rm e}^{\gamma t/t_{\rm ff}}\,, \quad \gamma = 0.5\,\alpha\,(t_{\rm ff}/t_{\rm d}) - \eta\,\epsilon_{\rm ff}\,f_{\rm g}\,. \tag{66}$$

With our fiducial values ($\alpha = 0.33$, $t_{\rm ff}/t_{\rm d} = 1/3$, $\eta = 1$, $\epsilon_{\rm ff} = 0.01$, $f_{\rm g} = 1$) we have $\gamma = 0.045$. With $t_{\rm mig} = 8t_{\rm d}$ the growth factor during migration becomes $M(t_{\rm mig})/M_{\rm c} \simeq 2.9$. It requires a very strong wind of $\eta \sim 5.5$ for the mass-loss to balance the accretion and leave the clump with a constant mass till depletion, which in this case may occur before the clump completes its migration. For a significant mass-loss in a migration time, $t_{\rm mig} \sim 24t_{\rm ff}$, γ in equation (66) has to be significantly smaller than $-1/24$. With the fiducial value of $\alpha = 0.33$, this requires that $\eta\epsilon_{\rm ff}$ would be larger than its fiducial value of 0.01 by an order of magnitude. Alternatively, γ could obtain such negative values if the effective α is negative, e.g. representing a case where mass-loss by tidal stripping overwhelms the mass gain by accretion. However, the reported significant clump growth in the simulations, where both accretion and tidal stripping are at play, indicates that the effective α is positive and close to the assumed fiducial value. We conclude that a net mass-loss in the clumps is very unlikely.

4.4 Other implications

The predictions listed above have a few immediate and interesting implications. If winds are relatively efficient, i.e. $\psi_{\rm ej} \sim 1$, then when all types of stellar feedback are taken into account one expects giant clumps to experience fairly significant steady winds. Purely momentum-driven feedback is expected to provide a mass-loading factor η of order unity, and adiabatic supernova feedback can boost it to η of a few. We emphasize that the significant outflows hold even though the radiative trapping is negligible, and even though the clumps do not experience explosive disruption in a dynamical time.

If the clumps were allowed to reach depletion, the depletion time would have been of the order of a significant fraction of 1 Gyr. However, the clumps are likely to complete their migration inwards at a shorter time. During this migration, the clumps accrete mass from the disc and merge with other clumps, roughly doubling their mass in one orbital time.

The fact that the time-scale for migration is typically shorter than the time-scales for star formation and depletion indicates that the clumps complete their migration while still gas rich, thus taking part in the overall 'wet' inflow within the disc (Forbes, Krumholz & Burkert 2012; Cacciato, Dekel & Genel 2012; Dekel et al. 2013). This 'wet' inflow has interesting implications, e.g. it naturally leads to a compact bulge (Dekel & Burkert, in preparation) and could feed the central black hole (Bournaud et al. 2011, 2012).

A question often raised is whether the outflows from clumps can be the driver of turbulence in the disc, the mechanism that maintains the Toomre instability at $Q \sim 1$. A necessary condition is that the power in the outflows is comparable to the turbulence dissipative loss. The outflow power from $N_{\rm c}$ clumps is

$$\dot{E}_{\rm w} \sim N_{\rm c} \dot{M}_{\rm w} V_{\rm w}{}^2 \,, \qquad (67)$$

with $\dot{M}_{\rm w}$ and $V_{\rm w}$ as predicted above. The turbulence is expected to decay on a disc dynamical time, so the dissipation rate is

$$\dot{E}_{\rm dis} \sim M_{\rm g} \sigma_{\rm d}^2 / t_{\rm d} \,. \qquad (68)$$

One can use from our analysis above $V_{\rm w} = \nu V_{\rm c}$, $\dot{M}_{\rm w} = \eta \dot{M}_*$, $\dot{M}_* = \epsilon_{\rm ff} M_{\rm g}/t_{\rm ff}$. For a clump contraction factor c, the dynamical times are related as $t_{\rm d} = c t_{\rm ff}$. For a $Q \sim 1$ disc and clumps, one can estimate that the internal clump velocity and the external disc velocity dispersion are comparable (e.g. Ceverino et al. 2012); they are related by

$$V_{\rm c}{}^2 / \sigma_{\rm d}^2 \sim (\pi/2)\, c \,. \qquad (69)$$

Then, the ratio of the rate of energy injection by clump winds to energy loss due to decay of turbulence becomes

$$\frac{\dot{E}_{\rm w}}{\dot{E}_{\rm dis}} \sim 6\, N_{\rm c,5}\, \eta\, \epsilon_{\rm ff\,-2}\, \nu_3^2\, c_3^2 \,. \qquad (70)$$

This seems to indicate that there is enough energy in the outflows to continuously stir up the disc. However, it is likely that a large fraction of the outflow energy will be ejected along the descending density gradient perpendicular to the disc and not injected into the interclump medium in the disc plane, thus making the contribution of outflows to the disc turbulence only secondary. The gravitational gain by the VDI-driven inflow along the potential gradient within the disc is a more likely source of energy for maintaining the disc turbulence (e.g. Bournaud et al. 2011; Forbes et al. 2012; Cacciato et al. 2012; Dekel et al., in preparation).

5 COMPARISON TO OBSERVED CLUMPS

5.1 Observed clumps

Table 1 lists pioneering estimates of the properties of seven giant clumps as observed in five $z \simeq 2.2$ star-forming disc galaxies (SFG) using adaptive optics spectroscopy focusing on Hα at the ESO Very Large Telescope as part of the Spectroscopic Imaging survey in the Near-infrared with SINFONI (SINS) survey (Förster Schreiber et al. 2009). These data are based on table 2 of Genzel et al. (2011), with slight revisions for the massive clumps in ZC406690 from table 3 of Newman et al. (2012). The five galaxies are selected to be massive discs of rotation velocities $\sim$250 km s^{-1}, dynamical masses of more than $10^{11}\,{\rm M}_\odot$ within the inner 10 kpc, and SFR$\sim$120–290 ${\rm M}_\odot\,{\rm yr}^{-1}$. They sample the upper end of the SFG population, and therefore the most massive giant clumps.

The galaxies BX482 and ZC406690 are large clumpy rotating discs with a prominent $\sim$5 kpc ring of clumps and star formation. D3a15504 is a large rotating disc with a central AGN. ZC782941 is a more compact rotating disc, showing an asymmetry due to a compact clump off the main body of the galaxy, potentially a minor merger. BX599 is a compact system with a high-velocity dispersion and a small $\sim$3 kpc rotating disc.

The most prominent clumps were identified from at least two different maps of Hα velocity channels. Clump 1 in Table 1 is the dominant clump (A) in BX482, part of an $\sim$5 kpc ring that includes three additional smaller clumps. Clump 2 is an average over the six off-centre clumps (A-F) in D3a15504, none of which is particularly dominant over the others. Clump 3, from ZC782941, is at the largest distance from the disc centre and the brightest in Hα, while this galaxy shows four additional clumps closer to the centre. Galaxy ZC406690 shows four clumps in an $\sim$5 kpc ring, of which three were studied in Genzel et al. (2011) and listed here. Clump 4 is ZC406690-C, the brightest in *I*-band Advanced Camera for Surveys (ACS), faint in Hα and shows an elongated shape. Clump 5 is ZC406690-A, the brightest in Hα and rather round, compact and isolated. Its SFR is high, its stellar population is young and it is gas rich. Clump 6 is ZC406690-B, the second in Hα brightness and rather faint in *I*-band ACS. Its stellar population is rather old, and it is relatively gas poor. Clump 7 is an exception, the whole centre of the compact galaxy BX599, namely a compact star-forming bulge.

The first group of rows in Table 1 refer to the clump structural properties. The second group of rows are the observed SFR and wind properties. The third group is quantities deduced from the observed quantities. The quantities marked by asterisks '*' are directly deduced from the observations.

As described in Genzel et al. (2011) and Newman et al. (2012), the quoted quantities are highly uncertain. They are limited by resolution and by modelling assumptions. For example, the formal errors quoted in table 3 of Newman et al. (2012) are about 100 per cent for some of the quantities characterizing the winds. These pioneering observations should therefore serve as preliminary indications only.

The intrinsic clump radius $R_{\rm c}$ was determined from the half-width at half maximum (HWHM) of a Gaussian fit to the appropriate velocity channel after subtracting in quadrature the HWHM of the instrumental resolution. Since the latter is typically 2 kpc, larger than the intrinsic clump radius, the estimated $R_{\rm c}$ is rather uncertain.

The clump characteristic velocity $V_{\rm c}$ is derived here from the kinematic measurements of velocity dispersion σ and rotation $V_{\rm rot}$ assuming Jeans equilibrium: $V_{\rm c}{}^2 = \beta(V_{\rm rot}^2 + c\,\sigma^2)$. The steep clump density profiles dictate $c \simeq 3.4$ (Genzel, private communication), and $\beta \simeq 1.17$ (Genzel et al. 2011). Then, the dynamical clump mass is derived from $M_{\rm c} = G^{-1} V_{\rm c}{}^2 R_{\rm c}$. This gives larger masses than derived in Genzel et al. (2011) using $c = 2$, the value appropriate for an isotropic isothermal sphere. Genzel et al. (2011) evaluated the clumps' gas mass from the measured SFR using an adopted version of the KS law. With the recent calibration at $z \sim 2$ using CO measurements (Tacconi et al. 2013), the Kennicutt relation is not very different than equation (7) with $\epsilon_{\rm ff} \simeq 0.01$, and the estimated gas mass using the KS law with the recent calibration is similar to the dynamical mass as derived here. Note, however, that the dynamical mass could be underestimated if the clump deviates from equilibrium due to strong inflows or outflows.

Table 1. Observed properties of giant clumps form Genzel et al. (2011). Quantities marked by asterisks '*' are deduced relatively directly from the observations, while the other quantities are computed from them.

Clump no.	1	2	3	4	5	6	7
Clump name	BX482-A	D3a15504-A-F	ZC782941-A	ZC406690-C	ZC406690-A	ZC406690-B	BX599
z*	2.3	2.4	2.2	2.2	2.2	2.2	2.3
$R_{\rm c}$ (kpc)* [a]	1.0	1.0	0.8	1.2	0.8	1.2	1.5
$V_{\rm c}$ (km s^{-1})* [b]	125	111	195	159	163	187	152
$M_{\rm c}$ ($10^9 {\rm M}_\odot$) [c]	3.6	2.8	6.9	6.9	4.8	9.5	7.8
$t_{\rm ff}$ (Myr) [d]	7.7	8.7	3.9	7.2	4.7	6.2	9.5
$\dot{M}_*$ (M$_\odot$ yr^{-1})* [e]	12	3.3	17	14	40	11	66
$V_{\rm w}$ (km s^{-1})* [f]	350	400	420	355	440	810	1000
$\dot{M}_{\rm w}$ (M$_\odot$ yr^{-1})* [g]	12	3.6	34	13	117	78	185
$\epsilon_{\rm ff,\,-2}$ [h]	2.5	1.0	1.0	1.5	3.8	0.7	**7.9**
ν [i]	2.8	3.6	2.2	2.2	2.7	4.3	**6.6**
η [j]	1.0	1.1	2.0	0.9	2.9	**7.1**	2.8
$\psi_{\rm w}$ [k]	2	3	4	2	**6**	**34**	**14**
$\Sigma_{\rm gas}$ (g cm^{-2}) [l]	0.26	0.20	0.77	0.34	0.54	0.48	0.25
$t_{\rm dep}$ (Myr) [m]	302	817	274	517	**62**	218	**63**
ϵ_* [n]	0.50	0.48	0.33	0.52	0.25	0.12	0.26
$t_{\rm mig}/t_{\rm dep}$ [o]	0.63	0.27	0.36	0.35	1.9	0.71	3.7
$\epsilon_{*,\rm mig}$ [p]	~0.50	0.24	0.23	0.35	–	~0.12	–
$M_{\rm c,mig}/M_{\rm c}$ [q]	~0.50	0.73	0.54	0.68	–	~0.12	–

[a] $R_{\rm c} = R_{\rm HWHM}$ after beam smearing (of HWHM $\simeq 2$ kpc) is subtracted in quadrature.
[b] ${V_{\rm c}}^2 = \beta\,(V_{\rm rot}^2 + c\,\sigma^2)$, $\beta = 1.17$, $c = 3.4$, assuming Jeans equilibrium.
[c] $M_{\rm c} = G^{-1}{V_{\rm c}}^2 R_{\rm c}$.
[d] $t_{\rm ff} = R_{\rm c}/V_{\rm c}$.
[e] $\dot{M}_* = L({\rm H}\alpha)/2.1\times10^{41}$ erg s^{-1} extinction corrected.
[f] $V_{\rm w} = \langle V_{\rm broad}\rangle - 2\,\sigma_{\rm broad}$.
[g] Average of two photodissociation case-B models. Clumps 5 and 6 are from Newman et al. (2012).
[h] $\epsilon_{\rm ff,-2} = 4.2\,V_{\rm c,2}^{-3}\,\dot{M}_{*,10}$ (equation 49).
[i] $\nu = V_{\rm w}/V_{\rm c}$.
[j] $\eta = \dot{M}_{\rm w}/\dot{M}_*$.
[k] $\psi_{\rm w} = \dot{M}_{\rm w}V_{\rm w}/(L/c)$.
[l] $\Sigma_{\rm gas} = M_{\rm g}/(\pi {R_{\rm c}}^2)$, $M_{\rm g} = {\epsilon_{\rm ff}}^{-1}\dot{M}_* t_{\rm ff}$.
[m] $t_{\rm dep} = 1000\,{\rm Myr}\,(1+\eta)_2^{-1}\,\epsilon_{\rm ff,-2}^{-1}\,t_{\rm ff,10}$ (equation 55).
[n] $\epsilon_* = (1+\eta)^{-1}$ (equation 54, relevant when $t_{\rm mig}/t_{\rm dep} > 1$).
[o] $t_{\rm mig}/t_{\rm dep} = 0.25\,(1+\eta)_2\,\epsilon_{\rm ff,\,-2}$ (equation 59).
[p] $\epsilon_{*,\rm mig} = 0.24\,\epsilon_{\rm ff,-2}$ (equation 60, relevant when $t_{\rm mig}/t_{\rm dep} < 1$).
[q] $M_{\rm c,mig}/M_{\rm c} = 1 - \eta\,\epsilon_{*,\rm mig}$ (equation 61, ignoring accretion).

The SFR is derived from the Hα luminosity corrected for extinction, with an uncertainty of about 30 per cent. The clump outflow velocity $V_{\rm w}$ is estimated from the maximum blue-shift and width of the broad emission component – this is the main pioneering discovery of Genzel et al. (2011). Its error is about 33 per cent. However, by adopting the maximum wind velocity as the characteristic wind velocity $V_{\rm w}$ one may overestimate some of the calculated wind properties. The mass outflow rate $\dot{M}_{\rm w}$ is taken as the average of two different crude estimates using two photodissociation case B models, described in appendix B of Genzel et al. (2011). Because of the elaborate modelling involved, and the different results obtained from the different models, this quantity is naturally highly uncertain, with an error of the order of 100 per cent. This is what makes the current results indicative only, not to be taken too strictly on a case by case basis.

5.2 Comparison of theory to observations

The first four clumps, addressed as 'typical' clumps, seem to be consistent with the fiducial case discussed above for stellar momentum-driven outflows. The SFR efficiency $\epsilon_{\rm ff,\,-2}$ is of order unity, the wind velocity is two to four times the clump velocity, $\nu \sim 3$, the mass-loading factor η is about unity and in one case $\sim$2, and the momentum injection–ejection factor $\psi_{\rm w}$ is 2–3 and in one case $\sim$4, as predicted by the theoretical momentum budget discussed above.

The fifth clump, ZC406690-A, is unusual in terms of its high SFR of 40 M$_\odot$ yr^{-1}, but is still marginally consistent with the fiducial momentum-driven wind case. It shows a marginally high SFR efficiency of $\epsilon_{\rm ff,\,-2} \simeq 3.8$. Its outflow is on the strong side, with $\eta \simeq 2.9$ and $\psi_{\rm w} \simeq 6$, but this is still marginally consistent with the fiducial case. However, it is different in the sense that its high SFR and low $t_{\rm ff}$ yield a short depletion time of $\sim$60 Myr, which is about half the migration time. This clump will complete its migration intact only because the mass gain by accretion from the disc is expected to be larger than the mass-loss by outflow.

The last two clumps seem to be extreme cases of strong outflows that are inconsistent with stellar feedback, even when the adiabatic supernova feedback is at its maximum and the ejection into the wind is efficient. Clump 6 is the most extreme case in terms of outflow. It is the most massive clump, $\sim10^{10}$M$_\odot$, its SFR efficiency is rather typical, $\sim$11 M$_\odot$ yr^{-1}, with the stellar population rather old, but its outflow is excessive, with $\eta \simeq 7$ and $\psi_{\rm w} \simeq 34$. Based on our estimates of the momentum budget, such an outflow cannot be driven by stellar feedback. Either it requires another driving mechanism, the observational estimates are severe overestimates, or the SFR we measure today is substantially smaller than it was when the outflow was launched.

The bulge clump 7 is an exception, representing a whole galaxy rather than a Toomre clump embedded in a disc. It has the highest SFR $\simeq 66\,\mathrm{M_{\odot}\,yr^{-1}}$. It shows an outflow with a moderately large $\eta \simeq 2.8$ but with a very high momentum injection efficiency of $\psi_{\rm w} \simeq 14$. As a result, its depletion time of $t_{\rm dep} \simeq 63$ Myr is only a quarter of its migration time. According to our momentum budget, such a high value of $\psi_{\rm w}$ is more than what stellar feedback can offer in clumps; again, one could avoid this problem if the SFR we measure today is lower than it was when the bulk of the outflowing material was launched.

Andrews & Thompson (2010) proposed a scenario where multiple scattering is possible when the gas surface density is above a threshold of $\Sigma_{\rm gas}Z \sim 1.1\,\mathrm{g\,cm^{-2}}$. If this was true, and if the gas surface density in the extreme clumps was sufficiently high, this could have provided a possible explanation for the extreme clumps. However, the simulations of Krumholz & Thompson (2013) show that multiple scattering does not occur even at a very high surface density. Furthermore, we note that the gas surface density in all the observed clumps is in the range $(0.2-0.8)\,\mathrm{g\,cm^{-2}}$, below the suggested threshold value given that the metallicity is comparable to and slightly lower than solar. There seems to be a marginal correlation between $\Sigma_{\rm gas}$ and $\psi_{\rm w}$, but it does not help us explain the extreme outflows in clumps 6 and 7.

6 CONCLUSION

We have analysed the outflows expected from star-forming giant clumps in high-z disc galaxies that undergo VDI. We evaluated the outflow properties based on the momentum budget, namely the efficiency of momentum injection into the ISM per unit SFR by a variety of stellar momentum and energy sources. We then estimated the lifetime of the clumps given their VDI-driven migration towards the disc centre and the associated growth of clump mass by accretion from the disc.

Our results can be summarized as follows.

(i) Most of the mass-loss is expected to occur through a steady wind over many tens of free-fall times, or several hundred Myr, rather than by an explosive disruption in one or a few free-fall times, less than ∼100 Myr.

(ii) Radiation-hydrodynamics simulations by Krumholz & Thompson (2012, 2013) provide a key input to the momentum budget that radiation trapping is negligible because it destabilizes the wind. This means that each photon can contribute to the wind momentum only once, and the radiative force is limited to about L/c. This calls into question other recent works that assume a very large trapping factor without self-consistently computing it.

(iii) All the direct sources of momentum taken together inject momentum into the ISM at a rate of about $2.5\,L/c$. This includes radiation pressure, protostellar winds, main-sequence winds and direct injection of momentum from supernovae.

(iv) The early adiabatic phases in expanding supernova-driven shells and main-sequence winds, if they operate at maximum efficiency, bring it up to a total force of $5\,L/c$ for typical gas densities in the clumps. An unknown fraction of this force is actually used to drive the wind, so this can serve as an upper limit.

(v) The resulting outflow mass-loading factor is of order unity. If the clumps were allowed to deplete their gas into stars and outflows standing alone, the depletion time-scale would have been a few disc orbital times, a significant fraction of a Gyr, ending with about half the original clump mass in stars.

(vi) However, the clump migration time to the disc centre due to the VDI is of the order of an orbital time, about 250 Myr, so the typical clumps are expected to complete their migration prior to depletion.

(vii) Furthermore, based on analytic estimates and simulations, the clumps are expected to double their mass in a disc orbital time by accretion from the disc and mergers with other clumps, which overwhelm the mass-loss by tidal stripping. This high rate of gravitational mass growth implies a net growth of clump mass in time and with decreasing radius despite the continuous massive outflows.

(viii) From the six disc clumps observed so far, five are consistent with the predictions for stellar-driven outflows.

(ix) One extreme case shows an outflow with an estimated mass-loading factor of 7 and a momentum injection rate of $34\,L/c$. This may indicate that the observed outflow in this case is an overestimate, which is not unlikely given the large uncertainties in the observed properties. Otherwise, this may hint to a stronger driving mechanism. One possible way to obtain higher efficiencies is if the supernovae explode in extremely low density environments generated by the other feedback mechanism. Another possibility is that this clump is just now ending its star formation, and therefore the present measured SFR is smaller than the value that prevailed at the time most of the outflow was launched.

We conclude that stellar feedback is expected to produce steady massive outflows from the high-z giant clumps, with mass-loading factors of order unity and momentum injection rate efficiencies of a few. This is consistent with five of the six to seven observed giant clumps where outflows were observed so far, with one or two exceptions in which the estimated outflows are apparently stronger. Despite the intense outflows, which indicate gas depletion times of several hundred Myr, the clumps are not expected to disrupt by this process. Instead, they are expected to migrate to the disc centre on a somewhat shorter time-scale, roughly a disc orbital time or about 250 Myr (Dekel et al. 2009; Ceverino et al. 2010), and during this process they are expected to more than double in mass by accretion from the disc. One therefore expects the population of in situ clumps to show systematic variations in their properties as a function of radius in the disc, in the form of declining mass, stellar age and metallicity, and increasing gas fraction and specific SFR (Ceverino et al. 2012; Mandelker et al., in preparation). One also expects that the clump migration, combined with the VDI-driven interclump gas in the discs, is an efficient mechanism for forming compact spheroids (Dekel et al. 2009; Ceverino et al. 2010; Dekel et al., in preparation), and providing fuel for black hole growth and AGN activity (Bournaud et al. 2011).

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