Radiation pressure-driven galactic winds from self-gravitating discs

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ABSTRACT

We study large-scale winds driven from uniformly bright self-gravitating discs radiating near the Eddington limit. We show that the ratio of the radiation pressure force to the gravitational force increases with height above the disc surface to a maximum of twice the value of the ratio at the disc surface. Thus, uniformly bright self-gravitating discs radiating at the Eddington limit are fundamentally unstable to driving large-scale winds. These results contrast with the spherically symmetric case, where super-Eddington luminosities are required for wind formation. We apply this theory to galactic winds from rapidly star-forming galaxies that approach the Eddington limit for dust. For hydrodynamically coupled gas and dust, we find that the asymptotic velocity of the wind is \(v_\infty \approx 1.5 v_{\text{rot}}\) and that \(v_\infty \propto \text{SFR}^{0.36}\), where \(v_{\text{rot}}\) is the disc rotation velocity and SFR is the star formation rate, both of which are in agreement with observations. However, these results of the model neglect the gravitational potential of the surrounding dark matter halo and a (potentially massive) old passive stellar bulge or an extended disc, which act to decrease \(v_\infty\). A more realistic treatment shows that the flow can either be unbound or bound, forming a ‘fountain flow’ with a typical turning time-scale of \(t_{\text{turn}} \sim 0.1–1\) Gyr, depending on the ratio of the mass and radius of the rapidly star-forming galactic disc relative to the total mass and break (or scale) radius of the dark matter halo or bulge. We provide quantitative criteria and scaling relations for assessing whether or not a rapidly star-forming galaxy of given properties can drive unbound flows via the mechanism described in this paper. Importantly, we note that because \(t_{\text{turn}}\) is longer than the star formation time-scale (gas mass/star formation rate) in the rapidly star-forming galaxies and ultraluminous infrared galaxies for which our theory is most applicable, if rapidly star-forming galaxies are selected as such, they may be observed to have strong outflows along the line of sight with a maximum velocity \(v_{\text{max}}\) comparable to \(\sim 1.5 v_{\text{rot}}\), even though their winds are eventually bound on large scales.

Key words: galaxies: formation – galaxies: haloes – intergalactic medium – galaxies: starburst.

1 INTRODUCTION

Galactic-scale winds are ubiquitous in rapidly star-forming galaxies in both the local and high-redshift Universe (Heckman, Armus & Miley 1990; Heckman et al. 2000; Pettini et al. 2001, 2002; Shapley et al. 2003; Rupke, Veilleux & Sanders 2005; Sawicki et al. 2008). They are important for determining the chemical evolution of galaxies and the mass–metallicity relation (Dekel & Silk 1986; Tremonti et al. 2004; Erb et al. 2006; Finlator & Davé 2008; Peeples & Shankar 2011), and as a primary source of metals in the intergalactic medium (IGM; e.g. Aguirre et al. 2001). Moreover, galactic winds are perhaps the most extreme manifestation of the feedback between star formation in a galaxy and its interstellar medium (ISM). This feedback mechanism is crucial for understanding galaxy formation and evolution over cosmic time (Springel & Hernquist 2003; Oppenheimer & Davé 2006, 2008; Davé, Oppenheimer & Sivanandam 2008; Oppenheimer et al. 2010).

The most well-developed model for galactic winds from rapidly star-forming galaxies is the supernova-driven model of Chevalier & Clegg (1985), which assumes that the energy from multiple stellar winds and core-collapse supernovae in the starburst is efficiently thermalized. The resulting hot flow drives gas out of the host, sweeping up the cool ISM (Heckman, Lehnert & Armus 1993; De Young & Heckman 1994; Strickland & Stevens 2000; Strickland et al. 2002; Fujita et al. 2009; Strickland & Heckman 2009). Although this model is successful in explaining the X-ray properties of rapidly star-forming galaxies, the recent observational results that galaxies with higher star formation rates

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(SFRs) accelerate the absorbing cold gas clouds to higher velocities \((v \propto \text{SFR}^{0.5})\) and that the wind velocity is correlated with the galaxy escape velocity may challenge the traditional hypothesis that the cool gas is accelerated by the ram pressure of the hot supernova-heated wind, whose X-ray emission temperature varies little with SFR, circular velocity and host galaxy mass, indicating a critical galaxy mass below which most of the hot wind escapes (e.g. Martin 1999). These new observations may instead favour momentum-driven or radiation pressure-driven models for the wind physics (e.g. Martin 2005; Murray, Quataert & Thompson 2005, hereafter MQT05; Weiner et al. 2009).

The model that galactic winds may be driven by momentum deposition provided by radiation pressure from the continuum absorption and scattering of starlight on dust grains was developed by MQT05. However, the conclusions of MQT05 are based on an assumed isothermal potential and spherical geometry, and are thus most appropriate for bright elliptical/spheroidal galaxies in formation. On the other hand, the theory of radiation-driven winds from accretion discs from the stellar to galactic scales has also been studied (e.g. Tajima & Fukue 1996, 1998; Proga, Stone & Drew 1998, 1999; Proga 2000, 2003), but none of these works considered radiation from self-gravitating discs.

In this paper, we answer the question of whether or not large-scale winds can be driven by radiation pressure from self-gravitating discs radiating near the Eddington limit. These considerations are motivated by the work of Thompson, Quataert & Murray (2005, hereafter TQM05), who argued that radiation pressure on dust is the dominant feedback mechanism in rapidly star-forming galaxies, and that in these systems star formation is Eddington limited. For simplicity, throughout this paper we assume that the disc is of uniform brightness and surface density. In Section 2, we show that such discs are fundamentally unstable to wind formation because the radiation pressure force dominates gravity in the vertical direction above the disc surface. This result is qualitatively different from the well-known case in spherical symmetry. In Section 2, we also discuss the applicability of this model to rapidly star-forming galaxies, and then calculate the terminal velocity of the wind along the disc pole and its dependence on both the SFR and galaxy escape velocity. In Section 3, we assess the importance of a spherical stellar bulge or an extended passive disc, and dark matter halo potential. In Section 4, we discuss the three-dimensional wind structure and estimate the total wind mass-loss rate. We discuss our findings and conclude in Section 5.

## 2 RADIATION-DRIVEN WINDS AND THE TERMINAL VELOCITY

We consider the idealized model problem of a disc with uniform brightness and total surface density: \(I(r \leq r_{\text{rad}}) = I_0 \) and \(\Sigma(r \leq R_0) = \Sigma\), where \(r_{\text{rad}}\) and \(R_0\) define the outer radius of the luminous and gravitating portion of the disc, respectively. The flux-mean opacity to absorption and scattering of photons is \(\kappa\). The gravitational force along the polar axis above the disc is

\[
f_{\text{grav}}(z) = -2\pi G \Sigma \int_0^{R_0} \frac{z r_c dr}{(r^2 + z^2)^{3/2}} = -2\pi G \Sigma \left( 1 - \frac{z}{\sqrt{r_c^2 + z^2}} \right),
\]

and the vertical radiation force along the pole is

\[
f_{\text{rad}}(z) = \frac{2\pi c I}{\Gamma_0} \int_0^{R_0} \frac{z r_c dr}{(r^2 + z^2)^{3/2}} = \frac{2\pi I}{c} \frac{r_{\text{rad}}}{\sqrt{r_{\text{rad}}^2 + z^2}},
\]

The extra factor of \(z/\sqrt{r^2 + z^2}\) in the radiation force integral compared with the gravitational force is the projection cosine factor \(\cos \theta\) such that \(d F_{\text{rad}} \propto \cos \theta \, d\Omega\), where \(\theta\) is the angle between the \(z\)-axis direction and the direction of the solid angle \(d\Omega\) at the disc surface (Rybicki & Lightman 1979, see also Proga et al. 1998; Tajima & Fukue 1998). Thus, the Eddington ratio along the pole \(\Gamma(z) = f_{\text{rad}}(z)/f_{\text{grav}}(z)\) as a function of height \(z\) is given by

\[
\Gamma(z) = \Gamma_0 \left( \frac{r_{\text{rad}}}{r_D} \right)^2 \left( \frac{z^2 + r_c^2 + z^2}{z^2 + \frac{r_{\text{rad}}^2}{r_D^2}} \right),
\]

where \(\Gamma_0 = \Gamma(z = 0) = k I_0/(\Gamma_0 G \Sigma)\) is the Eddington ratio at the disc centre. A disc at the Eddington limit (\(\Gamma_0 = 1\)) requires \(I_{\text{rad}} = 2\pi G \Sigma /\kappa\), or flux \(F_{\text{rad}} = 2\pi G \Sigma /\kappa\). If \(r_{\text{rad}}/r_D > 1/\sqrt{2} \approx 0.7\), then \(\Gamma(z)\) increases along the \(z\)-axis above the disc. In particular, for \(r_{\text{rad}} \approx r_D\), the radiation force becomes twice the gravitational force as \(z \to \infty\):

\[
\Gamma_\infty = \Gamma(\infty \to \infty) = 2.
\]

Because \(\Gamma(z)\) increases monotonically with \(z\), an infinitesimal displacement of a test particle in the vertical direction yields a net vertical acceleration, and the disc is thus unstable to wind formation. This result for discs is qualitatively different from the spherical case with a central point source where \(\Gamma = f_{\text{rad}}/f_{\text{grav}}\) is constant with radius.

In the more realistic case, we consider an exponential disc with surface brightness \(I = I_0 \exp(-r/r_{\text{rad}})\) and surface density \(\Sigma = \Sigma_0 \exp(-r/R_0)\). The disc has approximately uniform brightness for \(r < r_{\text{rad}}\) and uniform density for \(r < R_0\), while the brightness and mass distribution cut off for \(r \approx R_{\text{rad}}\) and \(r \approx R_0\), respectively. The vertical radiation and gravitational forces along the pole above the exponential disc can be calculated by multiplying the integrals in equations (1) and (2) by the exponential terms \(\exp(-r/R_0)\) and \(\exp(-r/R_{\text{rad}})\), respectively. As a result, in this case the Eddington ratio becomes a function of \(R_{\text{rad}}/R_0\). We obtain the Eddington ratio along the pole \(\Gamma(z \to \infty) = 2\Gamma_0(z = 0)\) if \(r_{\text{rad}} = R_0\), and \(\Gamma(z \to \infty) \leq \Gamma_0(z = 0)\) if \(R_0 \geq \sqrt{2} R_{\text{rad}}\). Therefore, the exponential disc gives a result similar to the uniform disc case in that \(\Gamma(z)\) increases monotonically only if \(R_0 < \sqrt{2} R_{\text{rad}} \approx 1.4 R_{\text{rad}}\), otherwise the disc is not promising to drive a wind because the gravitational force dominates the outward radiation pressure force. Then we take \(r_D = \sqrt{2} R_{\text{rad}}\) and refer to the region outside of the active and bright region \((r \geq \sqrt{2} R_{\text{rad}})\) as the ‘extended passive disc’, and we discuss it together with the effects of an old stellar bulge and a dark matter halo in Section 3.

Here, we adopt the uniform disc model. If we consider the motion of a test particle in the outflow, the velocity along the \(z\)-axis can be written as

\[
v(z) = \frac{v_0^2}{4\pi G \Sigma r_D} = r_D \Gamma_0 \arctan \left( \frac{z}{r_D} \right) - \left( 1 + \frac{z}{\sqrt{1 + z^2}} \right),
\]

where \(\hat{r} = r_{\text{rad}}/r_D\), \(\hat{z} = z/r_D\) and \(v_0\) is initial vertical velocity. The first term on the right-hand side of equation (5) is the ‘radiation potential’ along the pole, while the second term is the gravitational potential. The right-hand side is always positive if \(\Gamma_0 \geq 1\) and \(\hat{r} \Gamma_0 > 2/\pi \approx 0.64\), and thus the gas can be accelerated to infinity. On the other hand, an unbound outflow is still possible in the sub-Eddington case \(\Gamma_0 < 1\) if the initial velocity \(v_0\) is sufficiently large to escape the gravitational potential above the disc until reaching the critical height \(z^*\) where \(\Gamma(z^*) = 1\), because the gas is decelerated from its initial \(v_0\) until it reaches \(z^*\), it will be accelerated to infinity if \(v(z^*) \geq 0\). Using this constraint, we calculate the minimum \(v_0\) required to drive an unbound wind in the sub-Eddington case, and
the critical height \( z^* \) where the wind profile acceleration changes sign. Fig. 1 shows the results. We introduce a characteristic velocity \( v_c = \sqrt{4\pi G \Sigma_0} \). The calculation is applied for \( v_0 = (2h/2\pi)^{1/2} \) or \( \tau_0 \approx 0.64 \). In this case, the wind can be accelerated at infinity. The important point here is that for a disc with sub-Eddington brightness, and non-zero vertical velocity \( v_0 \), the gas can be first decelerated and then accelerated to infinity. For the typical random initial velocities of the gas in galaxies \( \rho v_0^2 \approx 2\pi G \Sigma^2 \), we have \( v_0/v_c \sim (h/2\pi)^{1/2} \), which is the gas thickness scale \( h = 0.1r \) has been assumed for the second equality (Downes & Solomon 1998). For this reason we expect that even somewhat sub-Eddington discs can drive outflows, as shown in Fig. 1. In contrast with the spherical case for which \( v^2 - v_0^2 \propto \tau_0 \), uniformly bright self-gravitating discs allow the gas above the disc to be accelerated for the case with \( \tau_0 = 1 \) or even for some sub-Eddington cases. In Section 4, we show the three-dimensional trajectories of wind particles launched from both Eddington and sub-Eddington discs.

We wish to apply this simple theory to rapidly star-forming galaxies, which may reach the Eddington limit for dust (TQM05). However, this application depends on the extent to which a disc-like collection of point sources of radiation (stars) may be treated as a uniformly bright disc. In the case that the dusty gas of rapidly star-forming galaxies, the approximation of a uniform brightness disc is likely valid. However, for \( \Sigma \lesssim 10^{-3} \, \text{g} \, \text{cm}^{-2} \) (for a disc optically thin to ultraviolet (UV) radiation of massive stars. Because the massive stars are point sources and dominate the galaxy’s bolometric luminosity, the approximation of a two-dimensional uniformly bright disc breaks down. In the intermediate regime \( 10^{-3} \lesssim \Sigma \lesssim 0.1 \, \text{g} \, \text{cm}^{-2} \) the disc is optically thick to UV radiation, but optically thin to the re-radiated FIR, the application of the simple uniformly bright disc model is tentative, as the fraction of the light that is absorbed and scattered in the gas depends on the distribution of the sources relative to the dusty gas, the inter-star spacing relative to the vertical scale of the dusty gas and the grain size distribution, which affect the overall scattering albedo. We save a detailed investigation of the intermediate case for a future work, but here note two important points: (1) even for low-\( \Lambda \) galaxies, the fraction of diffuse (potentially scattered) UV light may be substantial (Thilker et al. 2005) and (2) if a finite thickness disc is uniformly emissive and absorptive, then \( \kappa_{\nu, \text{FUV}} F_{\text{FIR}} \sim \kappa_{\nu, \text{FIR}} F_{\text{FIR}} \), where \( \tau_{\nu, \text{FIR}} \sim \kappa_{\nu, \text{FIR}} \Sigma \), and \( F_{\text{FIR}} \) and \( F_{\text{FIR}} \) are the total, UV and FIR fluxes, respectively. The radiation pressure force above the disc is thus dominated by UV emission from the ‘skin’ of the disc (the \( \tau_{\nu, \text{FIR}} \approx 1 \)). Therefore, if the total flux is equal to the Eddington flux – \( F_{\text{Edd}} = 2\pi G \Sigma_{0}/\kappa \), where \( \kappa \) is the flux-mean opacity – then the escaping UV radiation is also at Eddington: \( F_{\text{Edd}} \sim \tau_{\nu, \text{FIR}} \sim \kappa_{\nu, \text{FIR}} \Sigma_{0}/(\tau_{\nu, \text{FIR}} \kappa) \). For these reasons, the model of a uniformly bright disc may apply even when \( \Sigma < 0.1 \, \text{g} \, \text{cm}^{-2} \) and the medium is not optically thick to the re-radiated FIR.

To proceed with our application to rapidly star-forming galaxies, the characteristic velocity \( v_c \) is written as

\[
v_c = \sqrt{4\pi G \Sigma_0} = 500 \, \text{km} \, \text{s}^{-1} \, \Sigma_0^{1/2} \, r_{D,1kpc}^{-1/2},
\]

where we take \( \Sigma_0 = \Sigma_{1}/1 \, \text{g} \, \text{cm}^{-2} = \Sigma/(4800 \, \text{M}_\odot \, \text{pc}^{-2}) \) and \( r_{D,1kpc} = r_{D}/1 \, \text{kpc} \). Momentarily neglecting the importance of the surrounding dark matter halo, or a potentially massive old stellar bulge or an extended passive disc (which we evaluate in Section 3), from equation (5) with \( \Sigma_0 = 1 \) and \( \tau_{rad} \approx r_0 \), the asymptotic terminal velocity along the pole is

\[
v_\infty = v_c \sqrt{\pi/2 - \frac{1}{n}} \approx 380 \, \text{km} \, \text{s}^{-1} \, \Sigma_0^{1/2} \, r_{D,1kpc}^{-1/2}.
\]

This expression for \( v_\infty \) can be related to the SFR using the Schmidt law, which relates the star formation surface density and gas surface density in galactic discs: \( S\Sigma_{\text{FIR}} \propto S\Sigma_{\text{gas}} \) (Kennicutt 1998). We approximate \( \Sigma_{\text{gas}} \approx 0.5 \frac{S\Sigma_{\text{FIR}}}{D} \), where \( \frac{S\Sigma_{\text{FIR}}}{D} \) is the disc gas fraction. Since \( S\Sigma_{\text{FIR}} \approx S\Sigma_{\text{D}} \), we have \( v_\infty \approx \Sigma_0^{1/2} \, r_{D} \propto \Sigma_0^{1/2} \, D \propto \Sigma_{\text{SFR}}^{1/2} \, D \propto \frac{S\Sigma_{\text{gas}}}{D} \propto \frac{S\Sigma_{\text{FIR}}}{D} \)

\[
\approx 150 \, \text{km} \, \text{s}^{-1} \, f_{g,0.5}^{0.5} \left( \frac{S\Sigma_{\text{FIR}}}{50 \, \text{M}_\odot \, \text{yr}^{-1}} \right)^{0.36} \, r_{D,1kpc}^{-0.21},
\]

which is consistent with the observation \( v_\infty \propto S\Sigma_{\text{FIR}}^{0.35 \pm 0.06} \) in low-redshift ULIRGs (Martin 2005) and \( v_\infty \propto S\Sigma_{\text{gas}}^{0.3} \) for high-stellar-mass and high-SFR galaxies at redshift \( z \approx 1 \) (Weiner et al. 2009, but see fig. 17 from Chen et al. 2010). The observed scatter at a given SFR may be caused by different \( r_{D} \), \( f_{g} \), \( r_{rad} \), \( r_{D} \), \( v_0 \), bulge and dark matter halo mass, and the time dependence of the wind properties as the stellar population evolves (see Section 3). Since we only use a simplified uniform disc model to derive equations (7) and (8), a more definitive comparison with the data should await a model with more realistic distributions of surface density, opacity and brightness. Moreover, we can give the criterion for driving a galactic wind in terms of SFR or \( S\Sigma_{\text{FIR}} \). For an Eddington-limited disc, we have \( F_{\text{Edd}} \propto \Sigma_{\text{gas}} \propto \Sigma_{\text{SFR}}^{0.71} \propto \frac{S\Sigma_{\text{gas}}}{D}^{0.71} \propto \frac{S\Sigma_{\text{FIR}}}{D}^{0.71} \).

\[
F_{\text{Edd}} \simeq 5 \times 10^{11} f_{g,0.5}^{0.53} \left( \frac{S\Sigma_{\text{FIR}}}{50 \, \text{M}_\odot \, \text{yr}^{-1}} \right)^{0.71} \, L_{\odot} \, \text{kpc}^{-2},
\]

\[
\simeq 3 \times 10^{12} f_{g,0.5}^{0.43} \left( \frac{S\Sigma_{\text{FIR}}}{50 \, \text{M}_\odot \, \text{yr}^{-1}} \right)^{0.71} \, L_{\odot} \, \text{kpc}^{-2},
\]

These bounds on \( \Sigma \) depend linearly on the gas-to-dust ratio.
where we have again employed the Schmidt law and $\kappa = 10\,\xi, cm^2 g^{-1}$ is the flux-mean dust opacity. Keep in mind that the dust opacity in the FIR limit is about $\kappa_{\text{FIR}} \sim 1-10 cm^2 g^{-1}$.

If we assume that the galactic disc is in radial centrifugal balance with a Keplerian velocity $v_{\text{rot}} \sim \sqrt{GM_\text{halo}/R}$, and we take the typical terminal velocity from equation (7) to compare it to the rotation velocity at the radius $r_0$ that $v_{\text{int}} = v_{\text{int}}(r = r_0)$, we have $v_{\infty} \approx 2\sqrt{\pi}/2 - 1 \approx 1.5 \, v_{\text{rot}}$.

For an exponential disc with height $H$, the peak rotation velocity is $v_{\text{peak}} \approx 0.5 \sqrt{GM_\text{halo} r_{\text{peak}}/R}$, which is approximately reached at $z \approx 1.3 \, R_{\text{peak}}$. This estimate is again made in the absence of a galactic bulge or an extended passive disc and a dark matter halo. That the terminal velocity of the wind increases linearly with the galactic rotation velocity is also consistent with the observational results in Martin (2005, fig. 7). Because the flat part of the rotation curve $v_r$ is typically comparable to the rotation velocity at the edge of the galactic disc (Sharma, Nath & Shchekinov 2011), we also have $v_{\infty} \approx 1.5 \, v_r$. However, the factor '1.5' in equation (10) cannot be strictly obtained due to the gravitational potential of the dark matter halo or a spherical old stellar bulge or an extended passive disc, which decreases the factor 1.5 and can cause the flow to become bound (see Section 3).

In their cosmological simulations of structure formation and IGM enrichment by galactic winds, Oppenheimer & Davé (2006) assumed a wind launch velocity $v_{\text{launch}} = 3\sigma \sqrt{GM_\text{halo}/R_{\text{vir}}}$, which is the factor 3 is an assumption (see also Oppenheimer & Davé 2008; Oppenheimer et al. 2010). Taking the relation between velocity dispersion and rotation velocity $\sigma = v_{\text{rot}}/\sqrt{2}$ for the isothermal sphere (Binney & Tremaine 2008), the wind launch velocity can be written as $v_{\text{launch}} \approx 2\, v_{\text{rot}}$, similar to equation (10). However, besides the initial wind launch velocity, an additional kick is given to the wind particles to overcome the halo potential, implying that the scaling of wind velocity $v_{\text{in}}$ in the Oppenheimer & Davé model has a higher amplitude than what we obtain from just the effect of radiation pressure on dust discussed in this paper. Nevertheless, equations (7)–(10) are essential for developing an understanding of the maximum envelope of values for $v_{\infty}$ and its dependence on the observed properties of rapidly star-forming galaxies. More physics of $v_{\infty}$ and its $z$-dependence in the gravitational potential well of a spherical old stellar bulge or a dark matter halo is presented in Section 3.

Finally, we note that the characteristic time-scale for the wind to reach its asymptotic velocity is

$$t_c \sim \frac{\sqrt{r_0/(4G\Sigma)}}{3.5 \times 10^6 \Sigma_0^{1/2} D_\text{D, kpc} \, \text{yr}},$$

which can be compared with the time-scale of a bright star-forming disc: $t_s \sim M_\text{gal}/SFR$. Again employing the Schmidt Law, we find that

$$t_s \sim 2 \times 10^7 f_{\text{g,0}} \Sigma_0^{-0.4} \text{yr}$$

or $t_s/t_c \sim 0.02 \Sigma_0^{0.1} D_\text{D, kpc}^{0.1}$. Since $t_s \ll t_c$, the wind can be accelerated to $\sim v_{\infty}$ with a time-scale of $t_c$ before the gas supply is depleted by star formation in a time-scale of $t_s$. Note that gas recycling in the ISM will give an even longer $t_s$ than the simple estimate $M_\text{gal}/SFR$.

### 3 Extended Disc, Bulge and Dark Matter Halo

The galactic bulge, extended passive stellar or gaseous disc and dark matter halo are important to the wind dynamics. If we do not consider the luminosity from the galactic bulge and extended disc, the bulge, extended disc and halo only act to decrease the asymptotic wind velocity and may cause the wind to fall back to the disc as a 'fountain flow'. For the galactic bulge, we employ a truncated constant density sphere of mass $M_\text{bulge}$ and spherical radius $r_{\text{bulge}}$. The effect of an old passive extended stellar bulge or disc of gas can be neglected within a height scale $z \lesssim (\Sigma/\Sigma_{\text{ext}}) r_{\text{D}}$ with $\Sigma$ and $\Sigma_{\text{ext}}$ being the surface density of a nuclear disc ($r < \sqrt{2} r_{\text{rad}}$) and an extended star/gas disc ($r > \sqrt{2} r_{\text{rad}}$), respectively. On the other hand, its effect at large scales is similar to the effect of an old stellar bulge, but with the following replacements: $M_\text{bulge} \rightarrow M_{\text{ext}}$ and $r_{\text{bulge}} \rightarrow R$, where $R = \sqrt{r^2 + z^2}$ is the cylindrical distance to the halo centre and $M_{\text{ext}}$ is the total mass of an extended disc. Note that for normal spirals and field galaxies in the local Universe, their extended star/gas discs can be very large compared to the region of more rapid star formation. As discussed in Section 2, where we identified the requirement $R_0 < \sqrt{2} r_{\text{rad}}$ for wind formation, they are thus not promising to drive galactic winds via the mechanism discussed here. Instead, bulgeless galaxies without extended discs, and with very massive central concentrations of star formation – e.g. those ULRGs with very large nuclear surface density and SFR, but with small $\Sigma_{\text{ext}}$ – are most promising. For simplicity, we do not consider the importance of a passive disc further here. Also, we adopt the NFW potential (Navarro, Frenk & White 1996) to describe the dark matter halo distribution $\rho_\text{halo}(r) \propto R^{-3}$ for $r < r_\text{c}$, where $r_\text{c}$ is the scale radius. For simplicity, in this section we take the uniform disc model with $r_{\text{rad}} = R_0$. The Eddington limit $\Gamma_0 = 1$ including the dark matter halo becomes

$$\frac{\pi a^2}{c} \frac{1}{\Gamma_0} = 2\pi GM_\text{halo} \Sigma_0^2 \frac{1}{\Sigma_{\text{ext}}^2} \frac{1}{\Sigma_{\text{vir}}^2} = 2\pi GM_\text{halo} \Sigma_0^2 \left(1 + \frac{f_h}{2c}\right),$$

where $\Sigma_{\text{vir}} = \Sigma_{\text{halo}} / \Sigma_0$ for $r = r_{\text{vir}}$ is the virial radius, $f(\Sigma_{\text{vir}}) = \ln (1 + \Sigma_{\text{vir}} / \Sigma_{\text{gal}})$ and

$$f_h = \frac{M_{\text{halo}}}{2\pi r_0 \Sigma_{\text{halo}}^2 f(\Sigma_{\text{halo}})} - \frac{M_{\text{halo}}}{2M_{\text{disc}}^2} \left(\frac{r_D}{r_h}\right) \frac{1}{f(\Sigma_{\text{halo}})} .$$

We introduce the parameters $c^\prime = r_c / r_0$ and

$$f_0 = \frac{1}{2} \left(\frac{M_{\text{halo}}}{M_{\text{disc}}^2}\right) \left(\frac{r_D}{r_{\text{halo}}}\right) .$$

The parameters $f_0$ and $f_h$ measure the importance of the halo and bulge, respectively, in determining the dynamics of the flow. The asymptotic velocity in terms of these parameters is (see equation 5)

$$\frac{v_{\infty}^2 - v_{\text{rot}}^2}{v_{\infty}^2} = \Gamma_0 \left(1 + \frac{f_h}{2c}\right) \frac{\pi}{2} - (1 + f_h + f_h).$$

Note that for $f_0, f_h \rightarrow 0$, equations (16) and (7) are equivalent. We combine the effects of the galactic bulge and dark matter halo using the parameter $f_0 = f_h + f_h$. Assuming $\Gamma_0 = 1$, the condition for matter to be unbound is

$$f_0 < \left(\frac{v_{\text{int}}^2}{v_{\text{rot}}^2} + \frac{\pi}{2} - 1\right) \left(1 - \frac{\pi}{4c}\right)^{-1} .$$

More generally, taking $f_0$ as a parameter, the terminal velocity from equation (16) is less than the value from equation (7). For the purposes of an estimate, taking $c^\prime \sim 10$, the relation between $v_{\infty}$ and $v_{\text{rot}}$ becomes

$$v_{\infty} \approx 1.5(1 - 1.6 f_0)^{1/2} v_{\text{rot}},$$

$^2$The case of a bright spherical bulge was considered in MQT05.
which shows that the dark matter halo potential well is too deep for particles to escape for \( f_0 \gtrsim 0.6 \).

Quantitatively the halo, extended disc and bulge should have very similar effects in decreasing the outflow velocity. For simplicity, we can first neglect the bulge/disc and focus on similar effects in decreasing the outflow velocity. For simplicity, we can first neglect the bulge/disc and focus on similar effects in decreasing the outflow velocity. For simplicity, we can first neglect the bulge/disc and focus on similar effects in decreasing the outflow velocity.

In this case, the typical value of \( f_0 \) is given by

\[
 f_0 \sim 0.25 \left( \frac{M_{\text{halo}}/M_{\text{disc}}}{30} \right) \left( \frac{r_0/r_\text{D}}{1/30} \right) \left( \frac{2}{f(c_{\text{vir}})} \right),
\]

where we have taken the representative values of \( M_{\text{halo}}/M_{\text{disc}} \sim 10–100 \) (e.g. Leauthaud et al. 2012), \( c' = r_\text{D}/r_0 \sim 10–100 \) and \( f(c_{\text{vir}}) \sim 2 \) (for \( c_{\text{vir}} \sim 15 \); Macciò, Dutton & van den Bosch 2008).

Thus, instead of the naive estimate without the dark matter halo in equation (10), which yields \( v_{\text{esc}} \sim 1.5 v_\text{rot} \), we obtain \( v_{\text{esc}} \simeq 0.5 v_\text{esc} \). Even so, equation (18) makes it clear that the relation \( v_{\text{esc}} \simeq 1.5 v_\text{rot} \) from equation (10) cannot be strictly achieved for Eddington-limited rapidly star-forming galaxies by the mechanism discussed in this paper since all systems have dark matter haloes, not to mention the fact that the Oppenheimer & Davé model includes an additional kick to increase the wind terminal velocity. To the extent that ‘2’ in the expression of the Oppenheimer & Davé model \( v_{\text{esc}} \sim 2 v_\text{rot} \) or ‘\( \sim 3.5 \)’ with an extra kick is required to match observations and simulations, additional physics such as an effectively super-Eddington galaxy luminosity or supernova-driven hot flows would be necessary to add to the models here (e.g. Hopkins, Quataert & Murray 2011a, b; Murray, Ménard & Thompson 2011).

The quantitative scaling relations for bound and unbound outflows, encapsulated in the factors \( f_0 \) and \( c' \), are illustrated in Fig. 2. The left-hand panel shows the contours of the asymptotic kinetic energy \( (v_{\text{esc}}/r_\text{D})^2 \) as a function of \( f_0 \) and \( c' = r_\text{D}/r_0 \) assuming \( v_0 = 0 \) for unbound outflows. The right-hand panel of Fig. 2 shows the contours of \( \log(f_0) \) on the pole and the disc radius for bound ‘outflows’. According to equation (19), if all else being equal, a disc with higher ratio of \( M_{\text{halo}}/M_{\text{disc}} \) has larger \( f_0 \) and smaller \( z_{\text{turn}} \), and the flow has a shorter time-scale for reaching the turning point \( t_{\text{turn}} \sim z_{\text{turn}}/v_\text{c} \sim t_s(z_{\text{turn}}/r_D) \) (see equation 11). The right-hand panel of Fig. 2 shows that \( z_{\text{turn}}/r_D \) can be larger than \( \sim 1000 \), implying that the flow reaches many tens of kpc in height above the disc before falling back towards the host on a time-scale of \( \sim \) Gyr.

Importantly, even if \( f_0 \) is large enough so that the flow is bound, one may still observe an outgoing wind from a rapidly star-forming galaxy for two reasons. First, there is a region of parameter space where the time to reach the turning point \( t_{\text{turn}} \) is larger than the time for the rapidly star-forming galaxy to deplete its gas supply \( t_s = M_*/\text{SFR} \) (see equation 12). Since \( t_{\text{turn}}/t_s \sim (z_{\text{turn}}/r_D) t_s/t_\star \) (see equation 11), \( z_{\text{turn}}/r_D > 60 r_\text{D} v_\text{c}^2 f_0^{1/3} \Sigma_0^{0.4} \) is required for \( t_{\text{turn}}/t_\star > 1 \). Secondly, an observable flow along the line of sight to a non-edge-on disc can have a maximum velocity \( v_{\text{max}} \) that is comparable to \( v_{\text{esc}} \) in equation (7), even though the flow is bound on large scales \( (z_{\text{turn}}) \) by the dark matter halo. Fig. 3 shows the one-dimensional

![Figure 2. Contours of the asymptotic kinetic energy \((v_{\text{esc}}/r_\text{D})^2\) as a function of \( f_0 = f_h + f_b \) (assuming \( f_0 = 0 \)) and \( c' = r_\text{D}/r_0 \) for unbound particles (left), and the turning point \( \log(f_0(z_{\text{turn}}/r_\text{D})) \) along the pole in the case of bound particles, assuming \( v_0 = 0 \) (right). The dashed line in both panels divides ‘bound’ and ‘unbound’ trajectories.](https://academic.oup.com/mnras/article-abstract/424/2/1170/998627/)

![Figure 3. Outflow velocity \( v_\text{c}/v_\text{c} \) along the pole with \( c' = 10 \) and \( f_0 \) from top down \( f_0 = 0, 0.2, 0.4, 0.6, 0.8, 1.5 \) (thick lines), \( c' = 4 \) (dashed line) and \( c' = 100 \) (dotted line) with \( f_0 = 1 \). The horizontal line labelled as \( v_{\text{esc}}/v_\text{c} \) is the analytical solution as in equation (7).](https://academic.oup.com/mnras/article-abstract/424/2/1170/998627/)

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3 Increasing \( v_0 \) to 0.2\( v_\text{c} \) does not change the position of the contours appreciably.
outflow velocity $v_z/v_c$ along the pole with various $f_0$ and $c'$. In
the bound case, the maximum outflow velocities $v_{\text{max}}$ peak around
$1-10r_0$ before decreasing to much lower values. The maximum
$v_{\text{max}}$ is still comparable to $v_\infty$ in the absence of a halo or a bulge
or an extended disc, except for very large $f_0 \gtrsim 5$ (see equation 19).

The velocity of the flow before changing sign and falling back to
the host galaxy is time dependent and approximately reaches its
maximum at a time $\sim t_0$ (equation 11).

The fact that $t_\text{turn} > t_0$ and that $v_{\text{max}}$ approaches a few times $v_{\text{rot}}$
together imply that if galaxies are selected as bright rapidly star-
forming galaxies, they may appear to have unbound outflows even
though the gas is in fact bound on large scales by the dark matter
halo potential. Indeed, Fig.
3 implies that rapidly star-forming galax-
ies with bound flows may exhibit a rough correlation of the form
$v_{\text{max}} \sim 1-1.5 v_{\text{rot}}$.

### 4 THREE-DIMENSIONAL WINDS

AND MASS-LOSS RATE

So far we have focused on the forces along the polar $z$-axis. Fig.
4 shows the two-dimensional projections of the three-dimensional
orbits of test particles accelerated by radiation pressure and gravity,
as well as their constant time surfaces above a uniform
disc, starting from an initial height $z_0 = 0.1r_0$ and initial radii
$r_0 = 0, 0.1r_0, 0.2r_0, 0.3r_0, \ldots$, both with and without an NFW
potential, and with no galactic bulge or extended passive disc. The
opening angles of particle trajectories are formed because particles
initially corotate in the disc with an angular momentum $v_z^2/r$
except for particles along the $z$-axis. In the absence of a halo, the
upper-middle panel shows that winds can be driven from part of
the disc region (e.g. $r_0 > 0.5r_D$ in this panel) even when the disc
is sub-Eddington ($\Gamma_0 = 0.88$), while the upper-right panel shows
that the wind can also be launched even beyond the edge of a bright
disc ($r > r_D$) in a super-Eddington case ($\Gamma_0 = 5$). The lower pan-
els show that the character of the flow changes significantly as $f_0$
increases from 0.2, to 0.6, to 1.0 (left to right; compare with equa-
tion 19). For the case $f_0 = 0.2$, the wind is clearly unbound. For
$f_0 = 0.6$, particles emerging near the $z$-axis are accelerated to very
large vertical distances, whereas particles emerging from the outer
disc region fall back to the disc rapidly. A more massive halo with
$f_0 = 1.0$ produces only a ‘fountain flow’ in which particles fall back
to the disc with a time-scale of $\sim 0.11$ Gyr. Moreover, Fig.
5 gives the three-dimensional orbit vertical component $v_z$ evolution of the
particles driven from different disc regions. For a face-on disc, the
velocity $v_z$ of the wind is just the velocity along the line of sight.
The maximum velocities are reached roughly at $t \sim t_0$. Also, the
maximum and terminal velocities of particles from the outer disc
region are smaller than those from the inner disc region. As in Fig.
4, matter from the outer region of the disc can be bound by the halo
gravitational potential even though matter from the inner region is
unbound.

To estimate the mass ejection rate and mass-loss rate from
Eddington-limited discs, we first calculate that the Eddington flux

![Figure 4](https://academic.oup.com/mnras/article-abstract/424/2/1170/998627)

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and luminosity are

\[ F_{\text{Edd}} = 2 \pi c G \Sigma / \kappa \sim 3 \times 10^{12} \Omega_{\odot}^{-1} L_{\odot} \text{kpc}^{-2} \]  

and

\[ L_{\text{Edd}} = \pi \rho_{\text{rad}} F_{\text{Edd}} \sim 10^{13} \left( \frac{r_f}{D} \right)^2 \Omega_{\odot}^{-1} \Omega_{\text{kpc}} L_{\odot}. \]

respectively.

The total mass ejection rate from the disc surface \( M_{\text{ej}} \) is a local quantity measured on a disc scale height \( h_{\text{disc}} \). The disc, bulge and dark matter halo gravitational forces at the disc scale height \( h_{\text{disc}} \) are estimated at \( (R, z) \sim (0, h_{\text{disc}}) \) as \( f_{\text{disc}} \sim 2 \pi G \Sigma_{\text{disc}}, f_{\text{halo}} \sim \pi G \rho_{\text{halo}} h_{\text{disc}} \) and \( f_{\text{bulge}} \sim G M_{\text{halo}} / \left( 2 r_f^2 f(\nu_{\text{vir}}) \right) \), respectively. Taking the case of \( \rho_{\text{halo}} \rho_{\text{disc}} \ll \Sigma_{\text{disc}} \), the gravity of the bulge can be neglected at \( h_{\text{disc}} \) since \( f_{\text{bulge}} \ll f_{\text{disc}} \). Similarly, the ratio of the halo to the disc gravity \( f_{\text{halo}} / f_{\text{disc}} \sim 4 \times 10^{-3} \left( \frac{r_f}{D} \right) \left( \frac{M_{\text{disc}}}{M_{\text{halo}}} \right) \left( \frac{2}{f(\nu_{\text{vir}})} \right) \ll 1 \), the halo can also be neglected on the disc height scale \( h_{\text{disc}} \). Therefore, \( M_{\text{ej}} \) can be determined only by the disc and can be estimated in the absence of a bulge and a halo from the integrated momentum equation. In the single-scattering limit (i.e. all photons are scattered/absorbed once in the wind; e.g. MQT05), an estimate of \( M_{\text{ej}} \) is

\[ M_{\text{ej}} \sim 10^{13} \left( \frac{r_f}{D} \right)^2 \Omega_{\odot}^{-1} \Omega_{\text{kpc}} L_{\odot}, \]

where \( v_{\infty, f_0=0} \) is the terminal velocity of the flow without a bulge, an extended passive disc or a dark matter halo (equation 7). Since \( v_{\infty, f_0=0} \) is an upper limit to the velocity of the flow (valid as \( f_0 \to 0 \); see equation 18), equation (22) is only approximate. Nevertheless, to the extent that \( v_{\text{max}} \) is of the order of \( v_{\text{vir}} \) (see Fig. 3), equation (22) should yield an order-of-magnitude estimate of the mass-loss rate from the disc itself. Combining equations (7) and (22), we have

\[ M_{\text{ej}} \sim 3 \times 10^7 \left( \frac{r_f}{D} \right)^2 \Omega_{\odot}^{-1} \Omega_{\text{kpc}} L_{\odot} \text{yr}^{-1}, \]

which is similar to observational results (e.g. Martin 2005, 2006).

Using equations (8) and (9), the mass ejection rate is also given by

\[ M_{\text{ej}} \sim 6 \times 10^2 \left( \frac{f_{\text{SFR}}}{0.5} \right) \right)^{0.36} \left( \frac{50}{M_{\odot}} \right)^{0.30} \text{yr}^{-1}. \]

Using \( L = \epsilon \dot{c}^2 \text{SFR} \) to calculate the luminosity of a rapidly star-forming galaxy (e.g. Kennicutt 1998), where \( \epsilon_{-3} = \epsilon / 10^{-3} \) is the efficiency with which star formation converts mass into radiation, the single-scattering estimate for the ratio of the mass ejection rate to the SFR is

\[ \frac{M_{\text{ej}}}{\text{SFR}} \sim \frac{\epsilon c}{v_{\infty, f_0=0}} \sim 0.8 \epsilon_{-3} \Omega_{\odot}^{-1/2} \Omega_{\text{kpc}}. \]

which implies that radiation pressure drives more matter from discs in low-mass galaxies (see MQT05): for example, \( M_{\text{ej}}/\text{SFR} \geq 10 \) for \( \Sigma_{\text{disc}} \lesssim 3 \times 10^4 \Omega_{\odot} \Omega_{\text{kpc}}^{-1} \). However, as low-mass galaxies are usually the most dark matter halo dominated with higher ratio of \( M_{\text{halo}} / M_{\text{disc}} \) (Persic, Salucci & Stel 1996), or larger \( f_0 \) (see equations 17 and 19), the matter lost from low-mass discs may not escape the halo potential and will fall back on the time-scale \( t_{\text{fall}} \) (see the right-hand panel of Fig. 2). Otherwise, some additional physical mechanism beyond that described in this paper (e.g. a supernova-heated wind) is needed to unbind it from the halo.

Figure 5. Left panel: the time evolution of \( v_z/v_c \) for \( r / D = 0.2 - 1.0 \) starting with \( z / r_D = 0.1 \) and \( r_\parallel = 1 \). The horizontal line \( v_{\infty} / v_c \) is the same as Fig. 3. Right panel: the velocity \( v_z/v_c \) starting from \( (r / D, z / r_D) = (0.2, 0.1) \) (thick lines) and \( (0.8, 0.1) \) (thin lines) with different values of \( f_0 \) for the dark matter potential with \( c' = 10 \) and \( r_\parallel = 1 \).
However, both wind velocities and outflow rates can be significantly decreased by the presence of a spherical old stellar bulge or an extended gravitational disc or a dark matter halo potential (Section 3). Deeper or more extended spherical gravitational potentials cause the flow to be bound on large scales, and to produce only ‘fountain flows’ where particles fall back to the disc on a typical time-scale of $\sim 0.1$–$1$ Gyr, depending on the parameters of the system considered (see Figs 2–6). The criterion for the flow to become bound along the polar direction is given in equation (17). For typical values of the parameter $f_0$ (see equation 19), we find that the winds from rapidly star-forming galaxies can be either bound or unbound (see Fig. 2). As an example, for $f_0 \approx 0.25$ in equation (19), the asymptotic velocity of the wind is decreased by a factor of $\xi > 0.7$ from the case neglecting the dark matter halo completely (compare equations 10 and 18), from $\sim 1.5 v_{\text{rot}}$ to $\sim 0.5 v_{\text{rot}}$. However, for $f_0 > 0.6$ the asymptotic velocity goes to zero along the $z$-axis and the flow becomes bound on large scales. A more elaborate three-dimensional calculation shows that a totally unbound wind from a two-dimensional disc surface can only be approached at super-Eddington discs, otherwise the outflow asymptotic mass-loss rate from the large-scale gravitational potential $M_\infty$ is only a fraction of the mass ejection rate from the disc $M_0$ (Fig. 6).

Importantly, even in the limit of bound fountain flows, if the time-scale $t_{\text{run}}$ for reaching the turning point $z_{\text{run}}$ is longer than the lifetime of the rapidly star-forming galaxy $t_\star$ (equation 12), one may still observe an outward going wind while the rapidly star-forming galaxy is active and bright. The maximum positive velocity of the flow $v_{\text{max}}$, along the line of sight for non-edge-on discs is still correlated with the $v_{\text{rot}}$ of the disc (see Section 3). These facts may complicate the inference from observations of winds that they are unbound from the surrounding large-scale dark matter halo.

A more detailed work is required to fully assess radiation pressure on dust as the mechanism for launching a cool gas from rapidly star-forming discs. Extra observational evidence beyond the relation $v_{\infty} \propto \text{SFR}^{0.36}$ should certainly be added to verify the radiation-driven model from more realistic discs instead of the idealized model of uniform discs we considered in this paper. Equations (7) and (8) also give testable relations between $v_{\infty}$ and galactic surface density $\Sigma$, which can also be used to compare with observations of cold winds from galaxies with observed radii and masses. One key ingredient of our model is that galaxies should approach the dust Eddington limit. TQM05 proposed that a significant fraction of the radiation from ULIRGs should be produced by an Eddington-limited rapidly star-forming galaxy (see also the discussion of the dust Eddington limit in Hopkins et al. 2010). If so, the dust-driven mechanism discussed in this paper can be applied to the winds from ULIRG discs.

Moreover, galaxies at high redshifts are more active (e.g. Martin 2005, 2006; Weiner et al. 2009, see also Chen et al. 2010). The typical mass-loss rate from an Eddington-limited disc in the single-scattering limit is given by equation (25) and suggests that these outflows may efficiently remove mass from the disc (equation 25).
and (3) we ignore other physical mechanisms, which may act in concert with radiation pressure (MQT05; Murray et al. 2011). The general case with hydrodynamics, realistic disc brightness, surface density and dust opacity profiles will modify the picture presented here. Such an effort is underway. However, the basic conclusion that uniformly bright self-gravitating discs radiating near the Eddington limit are able to drive large-scale winds – particularly in the high-$\Sigma_1$ limit in rapidly star-forming galaxies (see Section 2) – should not be fundamentally changed by more elaborate considerations. Indeed, although we have specialized the discussion to rapidly star-forming galaxies and dust opacity, the instability derived in Section 2 is of general applicability.

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