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Acceleration of galactic supershells by \( \text{Ly} \alpha \) radiation

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ABSTRACT

Scattering of \( \text{Ly} \alpha \) photons by neutral hydrogen gas in a single outflowing ‘supershell’ around star-forming regions often explains the shape and offset of the observed \( \text{Ly} \alpha \) emission line from galaxies. We compute the radiation pressure that is exerted by this scattered \( \text{Ly} \alpha \) radiation on the outflowing material. We show that for reasonable physical parameters, \( \text{Ly} \alpha \) radiation pressure alone can accelerate supershells to velocities in the range \( v_{sh} = 200–400 \text{ km s}^{-1} \). These supershells possibly escape from the gravitational potential well of their host galaxies and contribute to the enrichment of the intergalactic medium. We compute the physical properties of expanding supershells that are likely to be present in a sample of known high-redshift \( (z = 2.7–5.0) \) galaxies, under the assumption that they are driven predominantly by \( \text{Ly} \alpha \) radiation pressure. We predict ranges of radii \( r_{sh} = 0.1–10 \text{ kpc}, \) ages \( t_{sh} = 1–100 \text{ Myr} \) and energies \( E_{sh} = 10^{53}–10^{55} \text{ erg} \), which are in reasonable agreement with the properties of local galactic supershells. Furthermore, we find that the radius, \( r_{sh} \), of a \( \text{Ly} \alpha \)-driven supershell of constant mass depends uniquely on the intrinsic \( \text{Ly} \alpha \) luminosity of the galaxy, \( L_{\alpha} \), the \( \text{H} \text{i} \) column density of the supershell, \( N_{\text{H}1} \), and the shell speed, \( v_{sh} \), through the scaling relation \( r_{sh} \propto L_{\alpha}/(N_{\text{H}1}v_{sh}^2) \). We derive mass outflow rates in supershells that reach \( \sim 10–100 \) per cent of the star formation rates of their host galaxies.

Key words: radiation mechanisms: general – radiative transfer – ISM: bubbles – galaxies: high redshift – cosmology: theory.

1 INTRODUCTION

The \( \text{Ly} \alpha \) emission line of galaxies is often redshifted relative to metal absorption lines (e.g. Pettini et al. 2001; Shapley et al. 2003). Scattering of \( \text{Ly} \alpha \) photons by neutral hydrogen atoms in an outflowing ‘supershell’ surrounding the star-forming regions can naturally explain this observation, as well as the typically observed asymmetric spectral shape of the \( \text{Ly} \alpha \) emission line (Lequeux et al. 1995; Lee & Ahn 1998; Tenorio-Tagle et al. 1999; Ahn & Lee 2002; Ahn, Lee & Lee 2003; Mas-Hesse et al. 2003; Ahn 2004; Verhamme, Schaerer & Maselli 2006; Verhamme et al. 2008). The presence of an outflow may also be required to avoid complete destruction\(^1\) of the \( \text{Ly} \alpha \) radiation by dust and to allow its escape from the host galaxies (Kunth et al. 1998; Atek et al. 2008; Hayes et al. 2008; Ostlin et al. 2008).

Indeed, the existence of outflowing thin shells (with a thickness much smaller than their radius) of neutral atomic hydrogen around \( \text{H} \alpha \) regions is confirmed by \( \text{H} \text{i} \) observations of our own (Heiles 1984) and other nearby galaxies (e.g. Ryder et al. 1995). The largest of these shells, so-called ‘supershells’, have radii of \( r_{max} \approx 1 \text{ kpc} \) (e.g. Ryder et al. 1995; McClure-Griffiths et al. 2002, 2006) and \( \text{H} \text{i} \) column densities in the range \( N_{\text{H}1} \sim 10^{19}–10^{21} \text{ cm}^{-2} \) (e.g. Lequeux et al. 1995; Kunth et al. 1998; Ahn 2004; Verhamme et al. 2008). The existence of larger extragalactic supershells has been inferred from spatially extended \( (\sim 1\text{~a few kpc}) \) \( \text{Ly} \alpha \) P-cygni profiles, that were observed around local star burst galaxies (Mas-Hesse et al. 2003). These ‘supershells’ differ from more energetic ‘superwinds’ (e.g. Heckman, Armus & Miley 1990; Martin 2005), in which galactic-scale biconical outflows break-out of the galaxies’ interstellar medium (ISM) with high velocities (up to \( \sim 10^3 \text{ km s}^{-1} \), of which M82 is a classical example). Supershells are thought to be generated by stellar winds or supernovae explosions which sweep-up gas into a thin expanding neutral shell (see e.g. Tenorio-Tagle & Bodenheimer 1988, for a review). The backscattering mechanism attributes both the redshift and asymmetry of the \( \text{Ly} \alpha \) line to the Doppler boost that \( \text{Ly} \alpha \) photons undergo as they scatter off the outflow on the far side of the galaxy back towards the observer.

In this paper, we show that the radiation pressure exerted by backscattered \( \text{Ly} \alpha \) radiation on the outflowing gas shell can

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\(^1\) \( \text{Ly} \alpha \) may also avoid complete destruction by interstellar dust when it is primarily locked up in cold clumps embedded in a hot medium that is transparent to \( \text{Ly} \alpha \) (Neufeld 1991; Hansen & Oh 2006).
accelerate it to velocities that reach hundreds of km s\(^{-1}\). Our basic result can be illustrated through an order-of-magnitude estimate. The net outward force exerted on a shell of gas\(^2\) by backscattered Ly\(\alpha\) radiation of luminosity \(L_{\text{BS,}\alpha}\) is \(d(m_\text{sh}v_\text{sh})/dt = L_{\text{BS,}\alpha}/c\). Here, \(m_\text{sh}\) is the total mass in the shell, \(v_\text{sh}\) is the shell velocity and \(c\) is the speed of light. If we assume that backscattering occurs over a time-scale \(\tau_{\text{BS}}\) and the shell mass is constant in time, then the net velocity gain by the shell is \(\Delta v_\text{sh} \sim 250\text{ km} \text{s}^{-1} (m_\text{sh}/10^7\text{ M}_\odot)^{-1} (\tau_{\text{BS}}/50\text{ Myr})(L_{\text{BS,}\alpha}/10^{33}\text{ erg s}^{-1})\). The adopted \(m_\text{sh}\) value corresponds to a thin spherical H\(\text{I}\) shell with a column density \(N_{\text{HI}} = 10^{20}\text{ cm}^{-2}\) and radius \(r = 1\text{ kpc}\). The fiducial values of \(L_{\text{BS,}\alpha}\) and \(\tau_{\text{BS}}\) were chosen to yield the typical observed shell speed and expected shell lifetime in Lyman-break galaxies (LBGs; see e.g. Shapley et al. 2003; Verhamme et al. 2008). Here, and throughout this paper, we assume that the supershells are mostly neutral. This assumption is based on 21-cm observations of local supershells which indicate that these are very thin (thickness \(<\) their radius, e.g. Heiles 1984) and therefore dense. At these high densities one expects free electrons and protons to recombine efficiently and make the gas significantly neutral. However, Pettini et al. (2000) argue that, based on the observed reddening of the ultraviolet (UV) continuum spectrum, either the dust-to-gas ratio in the supershell of LBG MS1 S12−cB58 is higher than the galactic value or alternatively that the H\(\text{I}\) only makes up \(f_{\text{HI}} \sim 14\sim 33\%\) of the total gas mass of the shell, with the remainder of the gas being molecular or ionized. We discuss how this latter possibility affects our results in Section 3.

This paper is a closely related to another paper in which detailed Monte Carlo Ly\(\alpha\) radiative transfer calculations are performed to assess Ly\(\alpha\) radiation pressure in a more general context (Dijkstra & Loeb 2008, hereafter Paper I). These radiative transfer calculations illustrate clearly that Ly\(\alpha\) radiation may be important in driving outflows in various environments. Importantly, the pressure exerted by Ly\(\alpha\) radiation alone can exceed the maximum possible pressure exerted by continuum radiation through dust opacity (see Paper I), which in turn can exceed the total kinetic pressure exerted by supernova ejecta (Murray, Quataert & Thompson 2005). This implies that Ly\(\alpha\) radiation pressure may thus in some cases provide the dominant source of pressure on neutral hydrogen gas in the ISM.

The goals of this paper are (i) to gauge the general importance of Ly\(\alpha\) radiation in the particular context of outflowing galactic supershells and (ii) to explore the related observable properties of Ly\(\alpha\)-driven supershells. In Section 2, we compute the time evolution of the velocity of an expanding shell of neutral gas that encloses a central Ly\(\alpha\) source. We consider a range of models, which are subsequently applied to known Ly\(\alpha\) emitting galaxies at high redshifts. Finally, we discuss our results and their implications in Section 3. The parameters for the background cosmology used throughout our discussion are \((\Omega_m, \Omega_\Lambda, \Omega_b, h) = (0.27, 0.73, 0.042, 0.70)\) (Dunkley et al. 2009; Komatsu et al. 2009).

\(^2\)The net momentum transfer rate per photon that enters the shell is \(\Delta p = (hv_\alpha/c)(1 – \mu)\), where \(\mu\) denotes a unit vector that is pointing radially outward. Furthermore, \(\mu = k_\text{LS} \cdot k_\text{out}\), where \(k_\text{LS}(k_\text{out})\) denotes the propagation direction of the Ly\(\alpha\) photon as it enters (leaves) the shell. When averaged over all directions, the total momentum transfer rate is \(\dot{N}_\alpha(hv_\alpha/c)e_\alpha\), where \(\dot{N}_\alpha\) is the rate at which Ly\(\alpha\) photons enter the shell. This total net momentum transfer rate applies regardless of the number of times each Ly\(\alpha\) photon scatters inside the shell or regardless of whether the Ly\(\alpha\) photon is absorbed by a dust grain (and possibly re-emitted in the infrared).

### 2 Pressure from Backscattered Ly\(\alpha\) Radiation

As mentioned in Section 1, the net momentum transfer rate from the Ly\(\alpha\) radiation field to the shell is given by

\[
\frac{d(m_\text{sh}v_\text{sh})}{dt} = m_\text{sh} \frac{dv_\text{sh}}{dt} = f_{\text{scat}}(v_\alpha, N_{\text{HI}}) \frac{L_\alpha}{c},
\]

where \(L_\alpha\) is the total Ly\(\alpha\) luminosity of the galaxy, and \(f_{\text{scat}}(v_\alpha, N_{\text{HI}})\) denotes the fraction of Ly\(\alpha\) photons that are scattered in the shell. The latter depends on various factors, including the column density of neutral hydrogen atoms in the shell, the intrinsic Ly\(\alpha\) spectrum (i.e. the emitted spectrum prior to scattering in the shell) as well as the shell velocity. For example, as \(v_\alpha\) increases, atoms in the shell interact with Ly\(\alpha\) photons that are farther in the wing of the line profile, and \(f_{\text{scat}}\) diminishes. Equation (1) assumes that the H\(\text{I}\) shell mass remains constant in time. This assumption is motivated by the observation that supershells are well-defined shells of H\(\text{I}\) gas that are physically separate from the central Ly\(\alpha\) source. These outflows differ from steady-state outflows (i.e. outflows in which mass is ejected at a constant rate \(\dot{M}\)) that occur around late-type stars (e.g. Salpeter 1974; Ivezic & Elitzur 1995). In reality, an H\(\text{I}\) supershell is not likely to maintain a constant mass as it sweeps through the ISM, but its evolution depends on the detailed hydrodynamics of the shell as it propagates through the ISM, which in turn depends on the assumed properties of the ISM. In this first study, we focus on simplified models that gauge the general importance of Ly\(\alpha\) radiation pressure on the dynamics of galactic supershells.

Furthermore, our calculation of the parameter \(f_{\text{scat}}\) conservatively ignores that possibility of ‘trapped’ Ly\(\alpha\) photons by the expanding supershell, which can boost the radiation pressure considerably. If radiation trapping were included, then the parameter \(f_{\text{scat}}\) would equal the ‘force-multiplier’ \(M_F\), which was shown to be much greater than unity at low \(v_\alpha\) in Paper I (see Section 3 for a more detailed discussion). We adopt a conservative approach because the calculation of \(M_F\) in Paper I assumed that the expanding shell contained no dust. Under realistic circumstances, the existence of dust inside the supershell will suppress the number of times the Ly\(\alpha\) photons can ‘bounce’ back and forth between opposite sites of the shell, which leads to a reduction in the value of \(M_F\). However, provided that the dust is located inside (or exterior to) the H\(\text{I}\) shell, the true value of \(M_F\) is always larger than the \(f_{\text{scat}}\) adopted in the paper (see Section 3).

#### 2.1 Calculation of \(f_{\text{scat}}\)

We compute \(f_{\text{scat}}(v_\alpha, N_{\text{HI}})\) under the conventional set of assumptions: (i) the emitted Ly\(\alpha\) spectrum prior to scattering is assumed to be a Gaussian with a velocity width \(\sigma = v_{\text{circ}}\) (Santos 2004; Dijkstra, Lidz & Wyithe 2007), where \(v_{\text{circ}}\) is the circular virial velocity of the host dark matter halo and (ii) the outflow is modelled as a single expanding thin shell of gas with a column density \(N_{\text{HI}}\) (as in Ahn et al. 2003; Ahn 2004; Verhamme et al. 2006, 2008). We can then write

\[
f_{\text{scat}}(v_\alpha, N_{\text{HI}}) = \frac{1}{\sqrt{2\pi\sigma^2}} \int dx \ e^{-x^2/2\sigma^2} \ e^{-N_{\text{HI}}\sigma_x(x-x_{\text{sh}})},
\]

where frequency is denoted by the dimensionless variable \(x \equiv (v – v_0)/\Delta v_{\alpha}\), with \(\Delta v_{\alpha} \equiv v_{\text{sh}}v_\alpha/c\) and \(v_\alpha\) being the thermal velocity of the hydrogen atoms in the gas given by \(v_\alpha = \sqrt{2k_B T/m_\alpha}\). \(k_B\) is the Boltzmann constant, \(T = 10^4\text{ K}\) is the gas temperature, \(m_\alpha\) is the proton mass, \(v_0 = 2.47 \times 10^4\text{ Hz}\) is the Ly\(\alpha\) resonance frequency, \(\sigma_x = \sigma/v_{\text{sh}}\), and \(x_{\text{sh}} = v_{\text{sh}}/v_\alpha\). In Fig. 1, we show \(f_{\text{scat}}\) as a function
The fraction of scattered Lyα photons $f_{\text{scat}}$ as a function of shell velocity $v_{\text{sh}}$ for a range of column densities $N_{\text{HI}}$ (in cm$^{-2}$, for $\sigma = 75$ km s$^{-1}$) and emitted line widths quantified by $\sigma$ (labelled by 75, 250 and 500 km s$^{-1}$ for log $N_{\text{HI}} = 20$). Higher $N_{\text{HI}}$ values require larger shell speeds in order for the Lyα photons to propagate through the shell unscattered because as $N_{\text{HI}}$ increases the photons need to be farther in the wing of the line profile (in the frame of the shell) for them not to scatter. Furthermore, a larger intrinsic linewidth, $\sigma$, serves to flatten the dependence of $f_{\text{scat}}$ on the shell speed (see text).

of $v_{\text{sh}}$ for a range of column densities $N_{\text{HI}}$ (for $\sigma = 75$ km s$^{-1}$) and emitted line widths $\sigma$ (for $N_{\text{HI}} = 10^{20}$ cm$^{-2}$).

As illustrated in Fig. 1, $f_{\text{scat}}$ increases with increasing $N_{\text{HI}}$, because photons need to be farther in the wing of the line profile (in the frame of the shell) for them not to scatter. For example, Fig. 1 shows that $\sim$50% of the Lyα photons are scattered when the shell speed is $\sim$200 km s$^{-1}$ for $N_{\text{HI}} = 10^{19}$ cm$^{-2}$, whereas the same fraction is reached when the shell speed is $\sim$600 (2000) km s$^{-1}$ for $N_{\text{HI}} = 10^{20}$ ($N_{\text{HI}} = 10^{21}$) cm$^{-2}$. Furthermore, increasing the intrinsic linewidth serves to flatten the dependence of $f_{\text{scat}}$ on the shell speed because at large $\sigma$ and low shell speeds there is a substantial fraction of photons far in the wing of the line profile.

### 2.2 The time evolution of the shell speed: general calculations

Given $f_{\text{scat}}(v_{\text{sh}}, N_{\text{HI}})$, we next use equation (1) to compute the evolution of the shell position ($r_{\text{sh}}$) and speed ($v_{\text{sh}}$) as follows.

1. The initial location and shell speed are denoted by $r_{\text{sh},0}$ and $v_{\text{sh}}$. If the outflowing shell covers $4\pi r_{\text{sh}}^2$ of the sky surrounding the Lyα emitting region (see Ahn 2004; Verhamme et al. 2008, for more detailed discussions), and if all the shell material is in neutral atomic hydrogen gas then $m_{\text{sh}} = 4\pi r_{\text{sh}}^2 N_{\text{HI}} m_p$, where $r$ is the radius of the shell, $m_p = 1.6 \times 10^{-24}$ g is the proton mass, and $N_{\text{HI}}$ is the initial column density of neutral hydrogen atoms.

2. We combine equations (1) and (2) to compute $d v_{\text{sh}}/d t$.

3. Over an infinitesimal time-step $\Delta t$ a new shell position and velocity are then generated as

$$r_{\text{sh}}(t + \Delta t) = r_{\text{sh}}(t) + v_{\text{sh}}(t) \Delta t$$
$$v_{\text{sh}}(t + \Delta t) = v_{\text{sh}}(t) + (d v_{\text{sh}}/d t) \Delta t.$$

At the new position, the shell’s column density reduces to $N_{\text{HI}} = N_{\text{HI}}(r_{\text{sh}}/r_{\text{sh},0})^{-2}$, and we return to step 2.

Our fiducial model has the following parameters. We adopt $\sigma = 150$ km s$^{-1}$, corresponding to the circular virial velocity of a dark matter halo of mass $M = 10^{11}$ M$\odot$ at $z = 5.7$ as appropriate for a typical host of known Lyα emitting galaxies (e.g. Dijkstra et al. 2007). We also adopt $L_\alpha = 10^{43}$ erg s$^{-1}$, corresponding to the typical observed luminosity of Lyα emitting galaxies (e.g. Ouchi et al. 2008). This Lyα luminosity corresponds to a star formation rate of $\sim 5 M_{\odot}$ yr$^{-1}$ for a Salpeter initial mass function and a gas metallicity $Z = 0.05 Z_{\odot}$ (Schaerer 2003, for $Z = 2 Z_{\odot}$ this star formation rate is higher by a factor of 2). For comparison, the total kinetic luminosity in supernova ejecta is $L_{\text{mech}} \sim 2 \times 10^{52} (E_{\text{mech}}/10^{51}$ erg) $N_{\text{SN}}/10^{22}$ erg s$^{-1}$, where $N_{\text{SN}}$ is the supernova rate per unit star formation rate and $E_{\text{mech}}$ is the total kinetic energy in ejecta per supernova explosion (e.g. Murray et al. 2005, and references therein).

For the shell, we assume $N_{\text{HI}} = 10^{20}$ cm$^{-2}$, $r_{\text{sh},0} = 1.0$ kpc (see Section 1 and Verhamme et al. 2008) and $v_{\text{sh},0} = 50$ km s$^{-1}$. We find that the final shell speed is not very sensitive to $r_{\text{sh},0}$. The above shell parameters imply a shell mass of $m_{\text{sh}} \sim 10^4 M_{\odot}$. The parameters of the fiducial model are summarized in Table 1.

The time evolution of the shell speed in the fiducial model is plotted as the solid line in the left-hand panel of Fig. 2, which shows that Lyα scattering accelerates the shell to $\sim 100$ km s$^{-1}$ after 10 Myr, consistently with the crude estimate given in Section 1. After 100 Myr, the shell speed reaches $\sim 230$ km s$^{-1}$. The acceleration of the shell decreases with time because $f_{\text{scat}}$ decreases with time. As the shell accelerates and expands, the fraction of Lyα photons that are scattered in the shell and the associated momentum transfer rate to the shell, decrease. Other lines in this panel represent models in which one of the model parameters was changed.

The red dashed line shows the shell evolution for a model in which the initial shell velocity was increased to $v_{\text{sh},0} = 150$ km s$^{-1}$. The difference in shell speed between this case and the fiducial model decreases with time. In the model represented by the green dot–dashed line the shell mass is reduced to $m_{\text{sh}} = 10^8 M_{\odot}$ by reducing $N_{\text{HI}}$. The shell speed evolves much faster, and reaches 200 km s$^{-1}$ after $\sim 5$ Myr. The final shell speed is $\sim 400$ km s$^{-1}$. The central (right) panel shows the shell velocity as a function of its radius (column density). In the fiducial model (and those in which one parameter was changed), the shell started out at $r = 1$ kpc with a relatively low column density. The model represented by the blue dotted lines is discussed in Section 2.3.

In all the above models, Lyα radiation pressure accelerates the shell up to a few hundreds of km s$^{-1}$. For comparison, the escape velocity from a halo with $v_{\text{esc}} = 150$ km s$^{-1}$ (which corresponds to a mass of $3 \times 10^{11} (1 + z)^{3/2} M_{\odot}$, e.g. Barkana & Loeb 2001) for a shell starting at $r = 1$ kpc is $v_{\text{esc}} = 440$ km s$^{-1}$ (at $z = 3$, with a very weak redshift dependence, see Section 2.3). Therefore, it is reasonable to claim that shells that are driven by Lyα radiation pressure may escape from massive dark matter haloes. Next, we consider whether this mechanism could operate in known observed high-redshift LBGs that have measured values of $L_\alpha$, $N_{\text{HI}}$ and $v_{\text{sh}}$.

### 2.3 The time evolution of the shell speed in known galaxies

In real galaxies, Lyα radiation is not the only process that determines the shell kinematics, and equation (1) modifies to

$$m_{\text{sh}} \frac{d v_{\text{sh}}}{d t} = f_{\text{scat}}(v_{\text{sh}}, N_{\text{HI}}) \frac{L_\alpha}{c} - F_{\text{mech}}(r) + 4 \pi r^2 \Delta P,$$

where $F_{\text{mech}}(r)$ is the mechanical energy input to the shell from the supernovae, and $\Delta P$ is the pressure difference between the inner and outer regions of the shell.

### Table 1. Parameters of Models used in Fig. 2.

| Number | $\sigma$ (km s$^{-1}$) | $L_\alpha$ (erg s$^{-1}$) | $r_{\text{sh},0}$ (kpc) | $v_{\text{sh},0}$ (km s$^{-1}$) | $N_{\text{HI},0}$ (cm$^{-2}$) |
|--------|---------------------|----------------------|-------------------|---------------------|------------------|
| 1      | 150                 | $10^{43}$            | 1.0               | 50                  | $10^{20}$         |
| 2      | 150                 | $10^{43}$            | 1.0               | 150                 | $10^{20}$         |
| 3      | 150                 | $10^{43}$            | 1.0               | 50                  | $10^{19}$         |
where $F_{\text{grav}}(r)$ is the gravitational force on the shell, and $\Delta P = P_{\text{int}} - P_{\text{ext}}$ with $P = \rho c^2$ being the pressure of the medium inside (‘int’) or outside (‘ext’) of the shell (Elmegreen & Chiang 1982). Here, $\rho$ is the density of the medium, and $c$ is its sound speed. The second and third terms on the right-hand side of equation (3) require assumptions about the unknown shape of the gravitational potential near the star-forming region and the local properties of the surrounding ISM.

To first order, gravity can be ignored in those galaxies for which the calculated terminal shell speed exceeds the escape velocity of the dark matter halo. More specifically, gravity can be ignored if the Ly$\alpha$ momentum transfer rate exceeds the gravitational force, $L_{\alpha}/c \gtrsim GM(<r_{\alpha})m_{\alpha}/r_{\alpha}^2$, where $M(<r_{\alpha})$ is the total mass enclosed by the supershell. Quantitatively, Ly$\alpha$ radiation pressure exceeds gravity when the Ly$\alpha$ luminosity exceeds

$$L_{\alpha} \gtrsim 10^{43} \text{erg s}^{-1} \left( M(<r_{\alpha}) \left[ \frac{10^8 M_\odot}{10^9 M_\odot} \right] \left( \frac{m_{\alpha}}{0.1 \text{ kpc}} \right)^{-2} \right).$$

If the radial density profile can be modelled as an isothermal sphere, then $M(<r_{\alpha}) \sim \rho_{\alpha}(r_{\alpha})r_{\alpha}^3$ and the required luminosity increases at smaller shell radii. However, if the shell starts at a small radius (i.e. $r_{\alpha,0} = 0.01 \text{ kpc}$), then for a fixed shell mass the column density is high (e.g. $N_{\text{HI},0} \gg 10^{22} \text{ cm}^{-2}$). This, in combination with the small initial velocity of the shell, implies that the Ly$\alpha$ photons are efficiently trapped and the impact of radiation pressure is bigger than estimated above (see Fig. 3 and Paper I). In Paper I, we showed that in this regime, a simple order-of-magnitude estimate for the importance of radiation pressure is obtained based on energy considerations (e.g. Cox 1985; Bithell 1990; Oh & Haiman 2002). Initially, the total gravitational binding energy of the shell is $|U_{\alpha}| = GM(<r_{\alpha})m_{\alpha}/r_{\alpha}$. The energy in the Ly$\alpha$ radiation field is $U_{\alpha} = L_{\alpha}t_{\alpha}$, where $t_{\alpha}$ is the typical trapping time of a Ly$\alpha$ photon in the neutral shell. The trapping time scales with the shell optical depth at the Ly$\alpha$ line centre, $\tau_0$, through the relation $t_{\alpha} \sim 15(\tau_0/10^5)r_{\alpha}/c$ (Adams 1975; Bonilha et al. 1979). Radiation pressure unbinds the gas when $U_{\alpha} / |U_{\alpha}| \gtrsim 1$. This is in agreement with the above discussion (except the first model of FDF 2002). Also $U_{\alpha} / |U_{\alpha}|$ for those galaxies in which the calculated terminal wind speed exceeds the escape velocity of the halo, which confirms that ignoring gravity is justified best for those galaxies in which the Ly$\alpha$ radiation pressure can accelerate the shell to a speed that exceeds the escape velocity of the host dark matter halo.

Since we investigate the properties of Ly$\alpha$-driven supershells, we assume that $\Delta P$ is less than the Ly$\alpha$ radiation pressure (this assumption is motivated by the possibility that Ly$\alpha$ radiation pressure provides the dominant source of pressure on neutral hydrogen gas in the ISM, see Section 1 and Paper I). Hence, we solve equation (3) under the assumption that the second and third terms on its right-hand side can be ignored. We initiate a shell at rest ($v_{\alpha,0} = 0$) at a very small radius ($r_{\alpha,0} = 0.01 \text{ kpc}$) and not at the origin to avoid a divergent $N_{\text{HI},0} \rightarrow \infty$. We get constraints on the mass, age and radius of the supershell by requiring that it reaches its observed expansion velocity at its observed column density (see Section 2.3.1 for an example).

Our results are summarized in Table 2. The first column contains the name of the galaxy. The second column contains the H$\text{\i}$...
column density and the third contains the shell speed as derived from fitting to the observed Lyα line profile (Schaerer & Verhamme 2008; Verhamme et al. 2008). The fourth column contains the intrinsic Lyα luminosity, which was inferred from the UV-based star formation rate, and the intrinsic Lyα equivalent width derived by Schaerer & Verhamme (2008) and references therein, which implies an approximate optical depth at 1216 Å of τ_{D} ∼ 1.3–3.3 (e.g. Verhamme et al. 2008). We argue in Section 3 that despite the presence of this dust, the Lyα radiation pressure can still be substantial (and possibly dominant over continuum radiation pressure).

2.3.1 MS1512–cB58

Pettini et al. (2002) have found that the lensed LBG MS1512–cB58 is surrounded by an outflowing shell of gas with a column density \( N_{H} = 7.5 \times 10^{20} \text{ cm}^{-2} \), which is expanding at a speed of \( v_{sh} = 200 \text{ km s}^{-1} \) (also see Schaerer & Verhamme 2008). The star formation rate inside the LBG, \( \sim 40 \text{ M}_{\odot} \text{ yr}^{-1} \), translates to an intrinsic Lyα luminosity of \( \sim 4 \times 10^{44} \text{ erg s}^{-1} \) (with the uncertainty due to the unknown metallicity of the gas, see Schaerer 2003). We adopt the central value of \( L_{\alpha} = 6 \times 10^{44} \text{ erg s}^{-1} \).

The blue dotted line in Fig. 2 depicts the time evolution of a shell with \( m_{sh} = 3 \times 10^{5} \text{ M}_{\odot} \) and the above parameters of MS1512–cB58. The plot shows that the observed shell velocity is reached at \( t \sim 1.9 \text{ Myr} \), when the shell reaches a radius \( r = 0.21 \text{ kpc} \) and has an observed column density of \( N_{H} = 7.5 \times 10^{20} \text{ cm}^{-2} \). More massive shells would reach the observed speed at later times, larger radii, and at a lower shell column density. The plot also shows that Lyα pressure accelerates the shell to \( \sim 360 \text{ km s}^{-1} \) after 10 Myr, at which point the column density has declined to \( \sim 4 \times 10^{18} \text{ cm}^{-2} \).

Lastly, we point out that the H I shell in MS1512–cB58 is known to contain dust, with an estimated extinction of \( E(B-V) = 0.3 \) (Schaerer & Verhamme 2008, and references therein), which implies an optical depth at 1216 Å of \( \tau_{D} \sim 1.3–3.3 \) (e.g. Verhamme et al. 2008). We argue in Section 3 that despite the presence of this dust, the Lyα radiation pressure can still be substantial (and possibly dominant over continuum radiation pressure).

2.3.2 Lyα emitters in the Focal Reducer and Spectograph deep field

Verhamme et al. (2008) reproduced the observed Lyα line profiles of 11 LBGs from the Focal Reducer and Spectograph (FORS) Deep Field at 2.8 < z < 5 observed by Tapken et al. (2007) based on a simple model in which the Lyα photons emitted by the LBGs backscatter off a single spherical outflowing shell with \( 2 \times 10^{19} \text{ cm}^{-2} < N_{H} < 7 \times 10^{20} \text{ cm}^{-2} \).

Table 2 contains two entries for FDF 1267 because two different models were found to reproduce the observed line profile. We have not included galaxies FDF 4691 and FDF 7539 in the table because these galaxies contained almost static supershells, which cannot be reproduced by our approach, possibly because the shells in these galaxies were slowed down by gravity.

Table 2. Predicted physical sizes of expanding H I supershells under the assumption that their expansion was driven predominantly by Lyα radiation pressure.

| Galaxy name | \( N_{H} \) (10^{20} \text{ cm}^{-2}) | \( v_{sh} \) (km s^{-1}) | \( L_{\alpha} \) (10^{43} \text{ erg s}^{-1}) | \( r_{sh} \) (kpc) | \( \theta_{sh} \) (arcsec) | \( m_{sh} \) (10^{4} \text{ M}_{\odot}) | \( t_{sh} \) (Myr) | \( E_{\alpha} \) (10^{54} \text{ erg}) | \( v_{esc} \) (km s^{-1}) | \( v_{term} \) (km s^{-1}) | f | M/SFR |
|-------------|-------------------------------|------------------|---------------------|----------|----------------|--------------------------|--------|-------------|------------------|------------------|-----------------|-------------|
| cB–58       | 7.5                           | 50               | 0.3                  | 5.5      | 0.09           | 58                       | 200    | 1.4         | 110.             | 544              | 170             | 1.0          | 0.24          |
| 1267        | 0.4                           | 150              | 5                    | 1        | 0.03           | 12                       | 7.4    | 1.0         | 164.             | 216              | 170             | 0.92         | 0.4           |
| 1337        | 0.2                           | 400              | 4.4                  | 0.03     | 0.04           | 0.67                      | 35     | 8           | 196              | 165              | 0.90            | 0.7          |               |
| 5550        | 0.2                           | 200              | 4.2                  | 0.3      | 0.03           | 2.1                       | 19     | 0.8         | 430              | 170              | 0.93            | 0.1          |               |
| 5812        | 5                             | 150              | 2.0                  | 4.3      | 0.63           | 35                       | 50     | 8           | 196              | 165              | 0.90            | 0.7          |               |
| 6557        | 0.2                           | 150              | 1.4                  | 0.26     | 11             | 2                         | 2.5    | 22          | 223              | 166              | 1.0             | 0.2          |               |
| 4691        | 0.8                           | 10               | 18                   | 230      | 4 × 10^{5}     | 4 × 10^{5}                | 67     | 170         | 10               | 1.0             | 1.9             | 0.3          |               |
| 7539        | 5.0                           | 25               | 42                   | 15       | 2.0            | 10^{4}                    | 10^{3} | 67          | 170              | 1.0             | 1.9             |             |               |

Note: FORS deep field Galaxies observed by Tapken et al. (2007). Parameters taken from Verhamme et al. (2008).

\[ f_{\text{esc}} = \frac{L_{\alpha}}{2\pi \mu_{\text{pc}} N_{\text{HI}} v_{\text{sh}}}. \]
For several galaxies, our model can be ruled out: our predicted $r_{sh}$ for FDF 4691 exceeds the size of the largest observed Ly$\alpha$ emitting structure in our Universe, which is $\sim 150$ kpc in diameter (Steidel et al. 2000). Furthermore, we find ages of $t_{sh} \sim 0.1$ Myr for two galaxies (FDF 5215 and FDF 1267). While this is not physically impossible, it appears unlikely that we have caught star-forming regions that early in their evolution. Indeed, these predicted ages fall well below the plausible age range that was derived for local supershells. Lastly, we predict a radius $r_{sh} = 15$ kpc for FDF 7531. Our model therefore predicts that existing observations should have resolved this galaxy as a spatially extended Ly$\alpha$ source (which is not observed).

For most other galaxies, $r_{sh} \sim 0.1$–1 kpc, $t_{sh} \sim 1$–200 Myr and $E_{sh} \sim 10^{53}$–10$^{55}$ erg. For comparison, McClure-Griffiths et al. (2002) find $r_{sh} \sim 0.07$–0.7 kpc, $t_{sh} \sim 0.9$–20 Myr and $E_{sh} = 0.03$–5 $\times 10^{33}$ erg for 19 galactic supershells. Our largest shells are larger than those of supershells that were observed in our galaxy, but are comparable in size of the inferred sizes of H$\alpha$ outflows around local star burst galaxies (Mas-Hesse et al. 2003). Thus, there is considerable overlap in the physical properties of observed (extra)galactic supershells and those in our model, although our model does contain a few supershells that are significantly larger and more energetic than observed galactic supershells. This may be because we ignore gravity. Indeed, for these energetic shells ($E_{sh} \sim$ a few $10^{54}$ erg), the maximum shell speed does not exceed the escape velocity of the host dark matter halo, and so gravity cannot be ignored.

We caution that the escape velocity of the dark matter halo that we computed is approximate. Our escape velocity was derived using the relation $v_{esc} = \sqrt{2v_{esc} \ln[f_{esc}/f_{esc,i}]}$, which assumes that the shell climbs out of the gravitational potential of an isothermal sphere. Although our calculated $v_{esc}$ depends only weakly on $r_{sh}$, we caution that deviations from our assumed simple model may change the results somewhat.

Interestingly, the fudge factor $f$ is within the range 0.65–1.0 in all models. The small scatter in $f$ arises because all shells initially have much higher H$\alpha$ column densities than their observed values. Therefore, for all shells the parameter $f_{scat} = 1$ during the early stages of their evolution. Depending on the precise time evolution of column density and shell speed, $f_{scat}$ eventually drops below unity (for some galaxies this has not happened yet and $f_{scat} = 1$, which in turn implies $f = 1$). Additional scatter may arise from the fact that some shells have predicted size that is close to the initial shell size assumed for our models.

The small scatter in $f$ implies that one can predict physical properties of galactic supershells without numerically integrating equation (3). Instead, one may simply adopt the central value $f \sim 0.9$ and apply equation (5) to quantities such as $L_u$, $N_{H\alpha}$ and $v_{sh}$ to predict $r_{sh}$. We also find that the mass outflow rates in supershells reach $\sim 10$–100 per cent of the star formation rates in their host galaxies, consistently with the total outflow rates that one expects theoretically (e.g. Erb 2008).

3 DISCUSSION AND CONCLUSIONS
It is well known that scattering of Ly$\alpha$ photons by neutral hydrogen in a single outflowing supershell around star-forming galaxies can naturally explain two observed phenomena: (i) the common shift of the Ly$\alpha$ emission line towards the red relative to other nebular recombination and metal absorption lines and (ii) the asymmetry of the Ly$\alpha$ line with emission extending well into its red wing. In this paper, we have computed the radiation pressure that is exerted by this scattered Ly$\alpha$ radiation on the outflowing shell.

We have shown that for reasonable shell parameters the shell can be accelerated to a few hundred km s$^{-1}$ after $\lesssim 10$ Myr. The Ly$\alpha$ acceleration mechanism becomes increasingly inefficient as the shell speed increases (Fig. 1) because the Doppler boost of the gas extends increasingly farther into the wing of the Ly$\alpha$ absorption profile allowing a decreasing fraction of the Ly$\alpha$ photons to scatter on the outflowing shell. Furthermore, for a shell mass that is constant in time, $N_{H\alpha} \propto r^{-3}$, and we find that for any choice of model parameters the shell achieves a terminal velocity that depends mainly$^5$. We found that equations (2) and (3) are accurate to within $\sim 50$ per cent on its total mass and the Ly$\alpha$ luminosity of the central source (Fig. 2).

We have computed the physical properties of expanding supershells that are likely to be present in specific observed high-redshift ($z = 2.7$–5.0) galaxies, under the assumptions that they are driven predominantly by Ly$\alpha$ radiation pressure and that their mass remains constant in time. We predict supershell radii that lie in the range $r_{sh} = 0.1$–10 kpc, ages in the range $t_{sh} = 1$–100 Myr and energies in the range $E_{sh} = 10^{53}$–10$^{55}$ erg, in broad agreement with the properties of local galactic supershells. We derive mass outflow rates in the supershells that reach $\sim 10$–100 per cent of the star formation rates in the host galaxies, in agreement with the total outflow rates that one expects theoretically (e.g. Erb 2008).

We have found that in all models the radius of the supershell is determined uniquely by parameters such as the intrinsic Ly$\alpha$ luminosity of the host galaxy (which may be inferred from the observed Ly$\alpha$ line shape), $L_u$, the total column density of H$\alpha$ in the supershell, $N_{H\alpha}$, and the shell velocity $v_{sh}$,

$$r_{sh} \sim 0.2 \left( \frac{f}{0.9} \right) \left( \frac{L_u}{10^{43} \text{erg s}^{-1}} \right) \times \left( \frac{10^{20} \text{cm}^{-2}}{N_{H\alpha}} \right) \left( \frac{200 \text{ km s}^{-1}}{v_{sh}} \right)^2. \quad (6)$$

Here, $f$ is a fudge factor that lies in the range $f = 1.0$–1.3. Our models have ignored gravity and the pressure of exerted by the surrounding ISM. This relation holds up only for supershells that are driven predominantly by Ly$\alpha$ radiation pressure. This relation is most accurate for those galaxies in which the shells are accelerated to velocities that exceed $v_{esc}$. We found this to be the case in five out of nine (55 per cent, see Table 2, ignoring the galaxies for which our model was ruled out, i.e. FDF 1267b, FDF 5215, FDF4691 and FDF 7539) of known galaxies. Hence, these shells could contribute to the enrichment of the intergalactic medium.

In other cases, gravity (and external pressure) will tend to reduce the value of the fudge factor $f$. Also, if neutral hydrogen gas only makes up a fraction $f_{H\alpha}$ of the total mass of the shell (see Section 1), then our predicted value of $r_{sh}$ should be lowered by a factor of $f_{H\alpha}$.

On the other hand, as was already mentioned in Section 2, our conservative calculations have completely ignored the fact that

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$^4$ The isothermal sphere model for the inner mass distribution of galaxies is supported observationally by gravitational lensing studies of elliptical galaxies (e.g. Winn, Rusin & Kochanek 2003; Bolton et al. 2008) and by the observed flatness of the rotation curves of spiral galaxies (e.g. Begeman 1989; Sanders & McGaugh 2002).

$^5$ An alternative way to see why the observed quantities $N_{H\alpha}$, $v_{sh}$ and $L_u$ provide unique constraints on the shell mass, radius and age is that these three unknown quantities relate to the observed quantities via three equations: (1) $m_{sh} = 4\pi r_{sh}^3 N_{H\alpha} \mu m_p$, (2) $v_{sh} \sim (L_u/\text{cm s}^{-1})/m_{sh}$ and (3) $r_{sh} \sim (L_u/2 \text{cm s}^{-1})/n_{H\alpha}$.
‘trapping’ of Lyα radiation by the optically thick supershell may significantly boost the Lyα radiation pressure compared to what was used in this paper. Specifically, it was shown in Paper I that when this trapping is properly accounted for, the total radiation force in equation (1) should include $M_F(v_{sh}, N_H)$ in place of $f_{rad}(v_{sh}, N_H)$. Here, $M_F$ is ‘force-multiplication factor’ that depends both the column density of H I and the velocity of the supershell. In Paper I, $M_F$ was computed for a range of $N_H$, and $v_{sh}$ values, by performing Monte Carlo Lyα radiative transfer calculation for each combination of $N_H$ and $v_{sh}$. In these calculations, the H I shell surrounds an empty cavity. The results of these calculations are shown in Fig. 3. As evident from the plot, $M_F$ greatly exceeds unity for low shell velocities and large H I column densities. Overplotted as the grey line is the trajectory in the $N_H-v_{sh}$ plane of the supershell in B–S (Section 2.3.1). At early times, $M_F > 10$, and we may have underestimated the radiation pressure force significantly (as was mentioned in Section 2, the actual force multiplication factor may be less than $M_F$, if gas and dust interior to the supershell prevents photons from repeatedly ‘bouncing’ back and forth between opposite sides of the H I shell). Lyα trapping could boost our predicted value of $v_{sh}$ by a factor as large as $(M_F)$, where $(M_F)$, the value of $M_F$, averaged over the expansion history of the shell.

So far, we have assumed H I to be the only source of opacity in the supershell. However, supershells may also contain dust which provides an additional source of opacity; for example, in the models of Verhamme et al. (2008), the optical depth through dust at $\lambda = 1216$ Å is $\tau_{dust} = 0.0$–2.0. At a given H I column density, we have therefore systematically underestimated $f_{rad}$, (absorption of a Lyα photon by dust also results in a momentum transfer of magnitude $h\nu/c$, which renders our calculations conservative. Furthermore, dusty supershells could also absorb significant amounts of continuum radiation, boosting the total momentum transfer rate even further. If the dust resides interior to the H I shell, then the Lyα flux that impinges upon the shell is reduced by $e^{-\tau_{dust}}$. In this case, the Lyα (as well as the continuum) radiation pressure is smaller than computed in this paper. However, the coupling between gas and dust makes it very unlikely that (Lyα) radiation pressure sweeps up the gas surrounding a star-forming region, while keeping the dust in place (see Murray et al. 2005, for more discussion on this).

Of course, the existence of dust inside the H I shell reduces its ability to ‘trap’ Lyα photons, which reduces the value of $M_F$ compared to that shown in Fig. 3. Despite this reduction, $M_F$ can still greatly exceed unity because at low shell expansion velocities the majority of Lyα photons do no penetrate deeply into the H I shell upon their first entry. Instead, the photons mostly scatter near the surface of the shell, and only ‘see’ a fraction of the dust opacity (indeed, for this reason Lyα radiation can escape from a dusty two-phased ISM, see Neufeld 1991; Hansen & Oh 2006). In other words, a relative small inner portion the H I shell can effectively trap Lyα photons. We have verified this statement with Monte Carlo radiative transfer calculations that included dust at the levels inferred by Verhamme et al. (2008): we typically found $M_F$ to be reduced by a factor of up to a few. In other words, even for dusty shells it is possible that $M_F \gg 1$.

Therefore, the pressure exerted by Lyα radiation alone can exceed the maximum possible pressure exerted by continuum radiation (also see Paper I), which can drive the large gas masses (as large as a few $10^{10}$–$10^{11}$ M⊙) out of galaxies in a galactic superwind (Murray et al. 2005). For comparison, in our model the Lyα photons transfer their momentum on to the expanding supershell, which contains significantly less mass ($m_{dust} \lesssim 10^8$ M⊙). This implies that in principle, both mechanisms could operate simultaneously, and that the neutral H I supershells may be accelerated to higher velocities than the bulk of the ejected gas. Since in our model Lyα radiation pressure operates only on a fraction of the gas, it does not provide a self-regulating mechanism for star formation and black hole growth as in Murray et al. (2005). However, it does provide a new way of enriching the intergalactic medium with metals.

We note that quasars can have Lyα luminosities$^7$ that reach $L_\alpha = 10^{46}$ erg s$^{-1}$ (Fan et al. 2006). In principle, this Lyα luminosity could transfer a significant amount of momentum on to neutral gas in its proximity. Because the Lyα emission line of quasars is typically very broad (∼few thousand km s$^{-1}$), $f_{rad}$ is significantly less than unity for the column densities considered in this paper (see Fig. 1). Furthermore, it is unclear whether gas clouds can remain neutral in close proximity to the quasar for an extended period of time. Even if this is the case, it remains to be determined whether the Lyα radiation pressure can compete with the continuum radiation pressure. The importance of Lyα radiation pressure near quasars therefore remains an open issue.

To conclude, observations indicate that it is the kinematics of H I gas surrounding star-forming regions that mostly determines the observed properties of the Lyα radiation (as opposed to dust content, e.g. Kunth et al. 1998; Atek et al. 2008; Hayes et al. 2008; Ostlin et al. 2008). This work suggests that – at least in some cases – the pressure exerted by the Lyα photons themselves may be important in determining the kinematics of H I gas. This is appealing for two reasons: (i) the mechanism operates irrespective of the dust content of the H I supershells. This mechanism may therefore operate also in galaxies of primordial composition at high redshift (especially since the total Lyα luminosity per unit star formation rate is higher, e.g. Schaerer 2003) and (ii) the shape and velocity offset of the observed Lyα emission strongly suggest that momentum transfer from Lyα photons to H I gas is actually observed to occur.

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$^6$ This term derives from the (time dependent) force-multiplication function $M(t)$ that was introduced by Castor, Abbott & Klein (1975), as $F_{rad} \equiv M(t)\tau_e\bar{v}_{sh}$/c. Here, $F_{rad}$ is the total force that radiation exerts on a medium, $\tau_e$ is the total optical depth to electron scattering through this medium. The function $M(t)$ arises because of the contribution of numerous metal absorption lines to the medium’s opacity, and can be as large as $M_{max}(t) \sim 10^3$ in the atmospheres of O-stars (Castor et al. 1975).

$^7$ Empirically, the Lyα luminosity of quasars is related to their B-band luminosity as $L_\alpha \sim \alpha L_B$ (Dijkstra & Wyithe 2006). High-redshift quasars have been observed with $L_B \gtrsim 10^{45}$ L⊙ = $4 \times 10^{45}$ erg s$^{-1}$ (see e.g. fig. 1 of Dijkstra & Wyithe 2006).

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