Galactic winds and the Lyα forest

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ABSTRACT
We study the effect of galactic outflows on the statistical properties of the Lyα forest and its correlation with galaxies. The winds are modelled as fully ionized spherical bubbles centred around the haloes in an $N$-body simulation of a $\Lambda$CDM model. The observed flux probability distribution and flux power spectrum limit the volume filling factor of bubbles to be less than 10 per cent. We have compared the mean flux as a function of distance from haloes with the Adelberger et al. (ASSP) measurement. For a model of bubbles of constant size surrounding the most massive haloes, bubble radii of $\gtrsim 1.5 h^{-1}$ Mpc are necessary to match the high transmissivity at separations $\lesssim 0.5 h^{-1}$ Mpc but the increase of the transmissivity at small scales is more gradual than observed. The cosmic variance error due to the finite number of galaxies in the sample increases rapidly with decreasing separation. At separations $\lesssim 0.5 h^{-1}$ Mpc our estimate of the cosmic variance error is $\Delta F \sim 0.3$, 30 per cent higher than that of ASSP. The difficulty in matching the rise in the transmissivity at separations smaller than the size of the fully ionized bubbles surrounding the haloes is caused by residual absorption of neutral hydrogen lying physically outside the bubbles but having a redshift position similar to the haloes. The flux level is thus sensitive to the amplitude of the coherent velocity shear near halos and to a smaller extent to the amplitude of thermal motions. We find that the velocity shear increases with halo mass in the simulation. A model where Lyman-break galaxies (LBGs) are starbursts in small-mass haloes matches the observations with smaller bubble radii than a model where massive haloes host the LBGs. If we account for the uncertainty in the redshift position of haloes, a starburst model with a bubble radius of $1 h^{-1}$ Mpc and a volume filling factor of 2 per cent is consistent with the ASSP measurements at the 1–1.5σ level. If this model is correct the sharp rise of the transmissivity at separations $\lesssim 0.5 h^{-1}$ Mpc in the ASSP sample is due to cosmic variance and is expected to become more moderate for a larger sample.

Key words: gravitation – intergalactic medium – cosmology: theory – dark matter.

1 INTRODUCTION

Stellar feedback, in the form of ionizing radiation and mechanical energy in galactic winds/outflows, plays an important role in determining the physical properties of galaxies and the intergalactic medium (IGM). The galactic contribution to the hydrogen-ionizing ultraviolet (UV) radiation may well exceed that of quasi-stellar objects (QSOs). Galactic superwinds regulate star formation, enrich the IGM with metals, and alter the dynamical and observational properties of the IGM close to galaxies.

Galactic winds are driven by the mechanical energy produced by supernova explosions of massive stars. They have been studied extensively in nearby starburst galaxies (e.g. Heckman, Armus & Miley 1990). The large velocity shifts between stellar and interstellar lines in the spectra of Lyman-break galaxies (LBGs) are strong evidence that galactic winds are also present in high-redshift galaxies (Pettini et al. 1998, 2001; Heckman et al. 2000). The inferred wind velocities are several hundred km s$^{-1}$, far above the sound speed of the IGM so that strong shocks are likely to be associated with these winds. These shocks will heat up, collisionally ionize, and sweep up the surrounding IGM near galaxies (e.g. Heckman 2001). Adelberger et al. (2003, hereafter ASSP) have analysed a sample of high-resolution QSO spectra with lines of sight (LOS) passing very close to LBGs. They detected a significant decrease of the Lyα absorption near galaxies. As discussed by ASSP and Kollmeier et al. (2003a,b) the increased UV flux close to LBGs falls short of explaining the decreased absorption by a factor of a few.

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decreased absorption is thus taken as evidence for the existence of dilute and highly ionized gas bubbles caused by wind-driven shocks.

Initially a shock propagating in the IGM expands adiabatically. During this phase the gas encountered by the shock is heated and collisionally ionized to become part of a bubble of dilute and hot plasma. When radiative losses become important, a thin cool shell forms at the shock front. The shell is driven outwards by the hot low-density inner plasma (e.g. Weaver et al. 1977; Ostriker & McKee 1988; Tegmark, Silk & Evrard 1993). Eventually the expansion of the plasma comes to a halt when pressure equilibrium with the surrounding IGM is reached. To explain the lack of HI close to the haloes, the bubble has to survive for a sufficiently long time. We defer the discussion of the conditions for the survival of the bubble to another paper. Here we assume that the bubbles are long-lived and study the implications of long-lived bubbles devoid of neutral hydrogen on QSO absorption spectra.

We will investigate how the correlation between the flux in the Lyα forest and the distance of the LOS to the closest galaxy depends on physical properties of the ionized bubbles around galaxies. For this purpose we use an \( N \)-body simulation of collisionless dark matter (DM). The basic observed properties of the Lyα forest can be reproduced with such simulations using a simple prescription that relates the distribution of gas and dark matter in the IGM (Hui & Gnedin 1997). With DM simulations it is possible to probe a wider dynamical range than with hydrodynamical simulations. This allows a better modelling of the statistical properties of the IGM–halo relation.

Measures of the halo–flux correlation, defined for example as the mean flux in the Lyα forest as a function of the distance to the nearest galaxies, are sensitive to the size of the bubbles. Absorption spectra probe the gas distribution in redshift space. Thermal motions and coherent peculiar velocities therefore strongly affect the halo–flux correlation at small separations (e.g. Kollmeier et al. 2003). Previous modelling of the impact of winds on the Lyα forest had difficulties in matching the large decrease of the absorption measured by ASSP at comoving separation \( s \leq 0.5 \, h^{-1} \) Mpc from the galaxies (Bruscoli et al. 2003; Croft et al. 2002; Theuns et al. 2002b; Kollmeier et al. 2003a,b). We will investigate here the role of peculiar velocities, thermal motions and the scaling of bubble size with halo mass in more detail. We will also perform a detailed assessment of the errors to obtain a more robust lower limit on the size and thus the volume filling factor of the bubbles.

The paper is organized as follows. In Section 2 we describe the model of the Lyα forest which we have used to calculate synthetic spectra from the DM simulations. In Section 3, we describe the simulation and the corresponding halo catalogues, compute three-dimensional flux maps of the simulation, and discuss the flux probability distribution function and power spectrum in the presence of winds. The main body of the paper is Sections 4 and 5 where we study in detail the impact of galactic winds on the halo–flux correlation and the corresponding observational errors. We conclude with a discussion of the results in Section 6.

2 THE LYMAN α FOREST AND THE MODERATE-DENSITY IGM

The Lyα forest is now widely believed to originate from an undiluting warm (10^4 K) photoionized IGM. The Lyα optical depth in redshift space due to resonant scattering can be expressed as a convolution of the real-space H I density along the line of sight with a Voigt profile \( \mathcal{H} \) (Bahcall & Salpeter 1965; Gunn & Peterson 1965):

\[
\tau(w) = \frac{C_0}{H(z)} \int_{-\infty}^{\infty} n_{\text{HI}}(x) \mathcal{H}[w-x-v_p(x), b(x)] \, dx, \tag{1}
\]

where \( C_0 = 4.45 \times 10^{-18} \) cm^2 is the effective cross-section for resonant line scattering and \( H(z) \) is the Hubble constant at redshift \( z \), \( x \) is the real-space coordinate, \( v_p(x) \) is the line-of-sight component of the H I peculiar velocity field, \( \mathcal{H} \) is the Voigt profile, and \( b(x) \) is the Doppler parameter due to thermal/turbulent broadening. For thermal broadening

\[
b(x) = 13 \left[ \frac{T(x)}{10^4 \, \text{K}} \right]^{1/2} \, \text{km s}^{-1},
\]

where \( T \) is the temperature of the gas.

Both \( x \) and \( w \) are measured in \( \text{km s}^{-1} \). The Voigt profile is normalized such that \( \int \mathcal{H} = 1 \) and can be approximated by a Gaussian for moderate optical depths,

\[
\mathcal{H} = \frac{1}{(\pi \beta^2)} \exp \left\{ -\frac{[\omega - x - v_p(x)]^2}{\beta^2} \right\}.
\]

The normalized flux obtained by ‘continuum fitting’ of the observed spectrum is related to the optical depth along the line of sight as

\[
F(w) \equiv \frac{I_{\text{obs}}(w)}{I_{\text{cont}}} = e^{-\tau(w)} \tag{2}
\]

where \( \tau \) is the optical depth, \( w \) is the redshift space coordinate along the line of sight, \( I_{\text{obs}} \) is the observed flux and \( I_{\text{cont}} \) is the flux emitted from the source (quasar) that would be observed in the absence of any intervening material.

On scales larger than a filtering scale that is related to the Jeans scale, \( x_J \), the IGM traces the DM distribution very well. For moderate overdensities the balance between adiabatic cooling and photoheating of the expanding IGM establishes power-law relations between temperature, neutral hydrogen density and total gas density (Katz, Weinberg & Hernquist 1996; Hui & Gnedin 1997; Theuns et al. 1998):

\[
T_g \propto \rho_g^\alpha \text{ and } n_{\text{HI}} \propto \rho_g^\beta,
\tag{3}
\]

where the parameter \( \beta \) is in the range 0.0–0.62, \( \alpha = 2 - 0.7 \beta \), and \( \rho_g \) and \( T_g \) are the density and temperature of the gas, respectively.

On scales smaller than the Jeans scale, pressure dominates over gravity. The gas pressure smoothing the gas distribution compared to the distribution of collisionless dark matter (e.g. Theuns, Schaye & Haehnelt 2000). We thus assume that \( \rho_{\text{HI}}(x) \propto \rho_{\text{dm}}^\gamma(x) \), where \( \rho_{\text{dm}}^\gamma(x) \) is the dark matter density smoothed in a way that mimics the pressure effects. We adopt the smoothing used in smoothed particle hydrodynamics (SPH) simulations (e.g. Springel, Yoshida & White 2001) to estimate the gas density from the dark matter particle distribution. The \( \text{H I} \) density \( n_{\text{HI}} \) and the smoothed dark matter field \( \rho_{\text{dm}} \) are then related as

\[
n_{\text{HI}}(x) = \hat{n}_{\text{HI}} \left[ \frac{\rho_{\text{dm}}^\gamma(x)}{\bar{\rho}_{\text{dm}}} \right],
\tag{4}
\]

where \( \rho_{\text{dm}} \) is the (volume) average of \( \rho_{\text{dm}}^\gamma \), and \( \hat{n}_{\text{HI}} \) is the \( \text{H I} \) density at \( \rho_{\text{dm}} = \bar{\rho}_{\text{dm}} \).

3 SIMULATING GALAXIES AND THE FLUX DISTRIBUTION

3.1 The DM simulation

One of our main goals is to assess the error in the observed galaxy flux correlation. This requires a simulation large enough to contain
a substantial number of DM haloes which can be identified as hosts of LBGs. At the same time the simulation should still reproduce typical absorption systems. This favours the use of a DM simulation for which a larger dynamical range can be achieved rather than a computationally more expensive hydrodynamical simulation. We use a $\Lambda$CDM simulation run with the $N$-body code GADGET (Springel et al. 2001). The cosmological parameters are listed in Table 1. It is part of a series of simulations of increasingly higher resolution which zoom in on a spherical region of initial comoving radius 26 $h^{-1}$ Mpc selected from a simulation with box size 479 $h^{-1}$ Mpc (cf. Stoehr et al. 2002). The inner region of the simulation thus has higher resolution than the outer regions. We only use the particle distribution inside a box of size $L = 30 h^{-1}$ Mpc which belongs to the spherical region of high resolution. The particle mass in this high-resolution region is $m = 9.52 \times 10^8 \, M_\odot/h$ (cf. Table 1). The large volume makes the simulation well suited for a direct comparison with the observed galaxy–flux correlation. We use a redshift output $z = 3$ to facilitate a comparison with the ASSP measurements.

### 3.2 DM Haloes as Hosts of LBGs

We aim at extracting halo catalogues with statistical properties similar to those of observed Lyman-break galaxies (LBGs). However, the relation between LBG and DM haloes is still somewhat uncertain. When LBGs were first discovered, a simple picture was advocated in which the mass of the DM haloes hosting the LBGs scales approximately linearly with the luminosity of LBGs, DM haloes are long-lived and each DM halo hosts one LBG (e.g. Steidel et al. 1996; Adelberger et al. 1998; Bagla 1998; Haehnelt, Natarajan & Rees 1998). In this ‘massive halo’ scenario the masses of the host haloes are rather large, $M \gtrsim 10^{12} \, M_\odot/h$, and the high-redshift LBGs are the progenitors of today’s massive and luminous galaxies. The rather slow decrease of the LBG space density with increasing redshift motivated an alternative picture, where most of the observed LBGs are interacting starburst galaxies in the process of assembling (e.g. Lowenthal et al. 1997; Trager et al. 1997; Kolatt et al. 1999; Somerville, Primack & Faber 2001). We will call this scenario the ‘starburst’ picture in the following. It predicts that most of the observed LBGs are hosted by haloes of smaller mass $M \sim 10^{11} \, M_\odot/h$, which will eventually merge to form more typical galaxies at $z \gtrsim 0$. Both scenarios predict similar clustering properties and cannot be ruled out based on the available clustering data (e.g. Wechsler et al. 2001). Direct observational constraints on the masses of LBGs are also still rather weak. Studies of rest-frame optical nebular emission lines imply halo masses in the range $10^{10}$–$10^{11} \, M_\odot/h$ (e.g. Pettini et al. 1998), while near-infrared imaging surveys combined with stellar population synthesis models and extinction yield virial masses of $M \sim 10^{11}$–$10^{12} \, M_\odot/h$ (e.g. Shapley et al. 2001). We will here first adopt the massive-halo picture, i.e. we assume that a dark matter halo contains only one LBG (e.g. Adelberger et al. 1998) and pick a minimum mass $M_{\text{min}}$ such that the number density of haloes $M \geq M_{\text{min}}$ in the simulation is similar to the observed number density of high-redshift LBGs. In Section 5.4 we will discuss what happens if we relax these assumptions.

DM haloes were identified with a friends-of-friends group-finding algorithm. Only groups containing at least 20 dark matter (DM) particles were classified as haloes. Since the DM particle mass is $9.52 \times 10^8 \, M_\odot/h$, the minimum halo mass is about $2 \times 10^{10} \, M_\odot/h$. Fig. 1 shows a slice of thickness $\Delta h = 0.21 \, h^{-1}$ Mpc at $z = 3$ in the simulation. The centres of the open circles show the location of haloes with $M \geq 5 \times 10^{10} \, M_\odot/h$. The radii of the circles are proportional to their masses.

In the simulation the number density of haloes with $M \geq 10^{11} \, M_\odot/h$ is about 0.04 $h^2$ Mpc$^{-3}$, whereas for haloes with $M \geq 10^{12} \, M_\odot/h$ it is about $2 \times 10^{-3} \, h^2$ Mpc$^{-3}$ (cf. Table 2). We have chosen $M_{\text{min}} = 5 \times 10^{10} \, M_\odot/h$ as our reference value but have also studied alternative choices. The corresponding number density is $4 \times 10^{-3} \, h^2$ Mpc$^{-3}$, which is about the number density of LBGs as determined for a $\Lambda$CDM cosmology at $z = 3$ from high-redshift surveys with limiting magnitude $25 \leq R \leq 27$ (Gaviláñez & Dickinson 2001). It should be noted that the halo cumulative function $n_s(m_h)$ as shown in Table 2 is consistent with a Sheth–Tormen mass function computed for a $\Lambda$CDM universe with identical parameters as our simulation (cf. Table 1).

### 3.3 Calculating synthetic spectra from the dark matter distribution

#### 3.3.1 From density to flux

The H I and the flux distributions are calculated from the DM distribution which was smoothed with a SPH kernel containing 20 particles within the smoothing length. The smoothed DM density
and velocity fields, $\rho_{\text{dm}}^i$ and $v_{\text{dm}}^i$ were interpolated for randomly chosen LOS through the simulation box. The pixel size of the spectra is $A = 60 \ h^{-1} \kpc$ and the comoving length $30 \ h^{-1} \Mpc$. From the density and velocity fields along the LOS we computed the optical depth as a function of the redshift space coordinate $w = x + v_p$, where $v_p$ is the component of $v_{\text{dm}}^i$ in the line-of-sight direction. Equation (1) then takes the following discrete form, \[
\tau(w_i) = A(z) \sum_j \left[ \rho_{\text{dm}}(x_j) \right]^\alpha \mathcal{H}_{ij},
\]
where $A = c \sigma_8 H(z)/(H(z)$ and we assume that $\rho_{\text{dm}}^i$ is normalized such that its mean is unity and $\mathcal{H}_{ij} = \mathcal{H}[w_i - x_j - v_p(x_j), b(x_j)]$ is a normalized Gaussian profile.

### 3.3.2 The temperature distribution

As we are using a DM simulation we still have to specify a temperature. For this we use the fact that temperature and density are tightly coupled. At moderate densities a simple power-law relation holds, $T \propto \rho_{\text{dm}}^\beta$ (cf. equation 3). In high-density regions, line cooling is efficient and the gas rapidly cools to $T \sim 10^4 \ K$ (e.g. Theuns et al. 1998; Davé et al. 1999). We will thus take the following relation between temperature and density

$$T_g = \begin{cases} T(\rho_{\text{dm}}^\beta) & \frac{\rho}{\rho_c} \leq 50 \\ 10^4 K & \frac{\rho}{\rho_c} > 50. \end{cases}$$

In the high-density regions where shock heating is important this is only a very crude approximation of the $T-\rho$ diagram of hydrodynamical simulation (Theuns et al. 1998; Davé et al. 1999; Croft et al. 2002). However, it is mainly the moderate-density regions which are responsible for the Lyman forest. Equation (6) should thus be sufficient for our purposes. The gas temperature $T$ at fixed overdensity is expected to evolve with redshift. Schaye et al. (2000) and Theuns et al. (2002a) claim to have detected a peak at $z = 3$ with decreasing temperature towards lower and higher redshifts. Note that we have ignored a possible redshift dependence and have assumed $T = 15000 \ K \simeq 10^4 \ K$ independent of redshift.

### 3.3.3 The effect of wind bubbles

We have adopted a very simple model for the effect of wind bubbles. We simply assume that shocks produce long-lived fully ionized spherical bubbles around galaxies (see Appendix A for a discussion of the physical properties of expanding wind bubbles). We assume that the neutral hydrogen $\text{H}^0$ density is zero in a spherical region of radius $r_w$ centred on the DM haloes which we have chosen to identify as LBGS. Note that this simple model has also been discussed in Croft et al. (2002), Kollmeier et al. (2003a, 2003b) and Weinberg et al. (2003). In the simulation, most of the haloes lie along filaments, and at the intersection of filaments, and the gas distribution around the haloes is often anisotropic on the scale of galactic winds (e.g. Croft et al. 2002; Springel & Hernquist 2003). Theuns et al. (2002) have also shown that winds propagate more easily into the low-density IGM than in the filaments, giving rise to highly ionized regions which are generally not spherical. It should thus be kept in mind that our assumption of a spherical wind bubble will only be a reasonable approximation for strong isotropic winds. Furthermore the bubbles do not necessarily live longer than a Hubble time. A discussion of the physical properties of shocks in the IGM can be found in Section A3 in Appendix A.

### 3.3.4 Properties of the synthetic spectra

Once we have smoothed the dark matter density $\rho_{\text{dm}}$, the flux distribution depends only on $A$ and $\alpha$ and the size distribution of wind bubbles. We have assumed $\alpha = 1.6$ and constrain $A$ from a random sample of 500 lines of sight of comoving length $L = 30 \ h^{-1} \Mpc$ by demanding that the mean flux $(F)$ is equal to the observed value $(F) = 0.67$ (Rauch et al. 1997; McDonald et al. 2000). The effect of varying the assumed size distribution of wind bubbles is discussed in the following sections. The spectral resolution is $\Delta \nu \simeq 6 \ h^{-1} \Mpc$, somewhat smaller than in the ASSP data, where it is $\Delta \nu \lesssim 10 \ km \ s^{-1}$. Note that the mean intercept distance in the simulation is $\sim 0.2 \ h^{-1} \Mpc$. This is sufficient to resolve $F$ on a scale $r \sim 0.25 \ h^{-1} \Mpc \sim 27 \ km \ s^{-1}$ as in the observations, but is somewhat too low to resolve all the features of the Lyman absorption lines (e.g. Theuns et al. 1998; Bryan et al. 1999). We also add noise with signal-to-noise ratio 50 per pixel of width $\Delta \nu = 2.5 \ km \ s^{-1}$.

### 3.4 Volume averaged statistics: The flux probability distribution function and the flux power spectrum

We will first investigate some basic flux statistics of our synthetic spectra to check if they are a reasonable representation of observed spectra and to study the effect of the wind bubbles. We begin with the flux probability distribution function (PDF) defined as the fraction of the volume of the IGM which has a flux value within a certain range. The right panel of Fig. 2 compares the pdf PDF at $z = 3$ for wind bubbles with different radii $r_w$. The filled circles show the observed PDF from McDonald et al. (2000), while the solid, long-dashed, short-dashed and dotted curves correspond to $r_w = 0$ (i.e. no wind), 0.5, 1 and 1.5 $h^{-1} \Mpc$ respectively. Following McDonald et al. (2000) we use 21 bins of width $\Delta F = 0.05$ to compute the flux PDF. The PDF of spectra without wind bubbles (solid curve) agrees very well with the PDF of the observed spectra (filled circles). Including wind bubbles with $r_w = 0.5 \ h^{-1} \Mpc$ has no significant effect on the flux PDF, and even when $r_w$ is as large as 1.5 $h^{-1} \Mpc$ there is only a slight discrepancy with the flux PDF of the observed spectra in saturated regions with $F \sim 0$. The small effect of the wind bubbles is due to their small volume filling factor $f_{\text{H}I}^w$, which is about 5 per cent for the model with $r_w = 1.5 \ h^{-1} \Mpc$. This is consistent with the findings of Theuns et al. (2002b) and Weinberg et al. (2003).

We can place an upper bound on the volume filling factor, $f_{\text{H}I}^w$, of bubbles by demanding a good match between the simulated and observed PDFs. We have computed the flux PDF for several bubble radii for which $f_{\text{H}I}^w$ is in the range 0.01–0.5. To account for the uncertainty in the mass of LBGs, we performed three computations for each $f_{\text{H}I}^w$: we selected DM haloes with mass $M$ larger than $5 \times 10^{10}$, $10^{11}$ and $5 \times 10^{11} \ M_{\odot}/h$, and adjusted $r_w$ which gives the desired $f_{\text{H}I}^w$. We find that independent of halo mass, a filling factor $f_{\text{H}I}^w \gtrsim 0.1$ yields a poor match to the observations, especially in the range $0.2 \lesssim F \lesssim 0.8$, as shown in the small plot inside the right panel of Fig. 2.

In the left panel of Fig. 2 we show the dimensionless one-dimensional flux power $\Delta F^2(k)$ as a function of $k$ (in $\text{km \ s}^{-1}$), defined as

$$\Delta F^2(k) = \frac{L_k}{\pi} |\delta_k|^2,$$

where $\delta_k$ is the Fourier transform of the flux contrast $(F/F - 1)$. The flux power spectrum was calculated with a fast Fourier transform (FFT) routine from 500 synthetic spectra. Note that our simulation box and therefore the spectra are not periodic which may
lead to small artefacts at small scales. The flux power of our synthetic spectra and that of the observed spectra obtained by McDonald et al. (2000) also agree very well. Bubbles of size as large as $r_w = 1.0$ (short-dashed line) and $r_w = 1.5 \, h^{-1}$ Mpc (dotted line). The results are compared to the observations of McDonald et al. (2000), shown as filled symbols. The triangles are the various volume averaged statistics in real space. The small plot in the right panel shows the PDF for filling factors of zero and ten per cent (cf. text). The wavenumber $k$ is in unit of km s$^{-1}$.

4 PROBING CORRELATIONS BETWEEN FLUX AND HALOES/LBGs

4.1 Statistical measures of the flux–halo correlation

There is a variety of statistical measures which can describe the correlation between the flux along a LOS and nearby galaxies. Here we will focus on the conditional flux function, $\bar{F}$, $\bar{F}$ is the mean flux at real-space distance $|s| = s$ to the next halo,

$$\bar{F}(s) = \frac{1}{N_h} \sum_{h} F(s | \text{halo}),$$

where $N_h$ is the total number of haloes, and $F(s | \text{halo})$ is the value of the flux at a distance $s$ from the nearest halo. Note that we use the position of the haloes which we identify as host haloes of LBGs as a proxy for the galaxy position. Results will thus depend on how this identification is done and we will later explore several possibilities. The summation is performed over a representative sample of haloes. The function $\bar{F}$ can be expressed in terms of the unconditional two-point correlation $\xi_{hl}$ as

$$\bar{F}(s) = \langle F \rangle \left[ 1 + \xi_{hl}(s) \right].$$

Relation (9) also holds in redshift space. In the following $\pi$ and $\sigma$ are the coordinates of $s$ parallel and perpendicular to the LOS, respectively.

4.2 Calculating the mean flux as a function of the distance to the nearest halo

For the numerical simulation we know the full three-dimensional flux field. We could therefore estimate $\bar{F}(s)$ from the transverse correlation $\bar{F}(\pi = 0, \sigma)$. However, real observation will sample the flux field $\bar{F}(s)$ more sparsely. ASSP have therefore estimated $\bar{F}(s)$ from the averaged 2D halo flux correlation function $\xi_w$ using equation (9) in order to maximize the signal. We proceed as ASSP do and calculate $\bar{F}(s)$ as follows. We first pick a LOS direction. We then produce an ensemble of spectra along this direction with random offsets perpendicular to the chosen direction. For each pair of halo/spectrum we determine the flux $\bar{F}(\pi, \sigma)$, the distance $\sigma$ at the point of closest approach, and the distance $\pi$ between the halo and a given pixel of the spectrum. This process is repeated for an ensemble of randomly chosen LOS directions. We then obtain $\bar{F}(s)$ by annular averaging $\bar{F}(\pi, \sigma)$, and binning in bins of comoving size $\Delta s$. We further average over all haloes and LOS. In the left panel of Fig. 3 we plot the full 2D distribution of the average flux $\bar{F}(\pi, \sigma)$ around haloes assumed to host LBGs. The haloes were assumed to have no galactic wind bubbles. The bin size is $0.25 \, h^{-1}$ Mpc, twice smaller than that we use for $\bar{F}$. The shaded area shows bins for which $\bar{F}$ is within $0.05$ of the mean flux $\bar{F}$. Contours are for flux levels increasing from 0.2 to 0.6 with decreasing line width. The contours are compressed along the line of sight as a result of the peculiar and thermal motions of the gas. At small separation the transmissivity is much smaller than the mean due to the increased density near haloes.

Redshift distortions cause $\bar{F}(s)$ (as computed in ASSP) to be different from the transverse correlation $\bar{F}(\pi = 0, \sigma)$. This can be seen in the middle panel of Fig. 3. The effect is important for $s \sim 1–3 \, h^{-1}$ Mpc. Note that previous work compared also the ASSP measurements with $\bar{F}(s)$ obtained from numerical simulations (e.g. Croft et al. 2002; Bruscoli et al. 2003; Kollmeier et al. 2003).

4.3 Estimating errors due to sparse sampling

The sample of ASSP contains 431 galaxies in seven fields containing a QSO (one field contains two QSOs). At large scales (> a few
function $N$ averaged over the halo/spectrum pairs lying in the given bin. The cumulative halo function $F(s)$ is only about 110. This means that we severely oversample the halo regions of angular dimensions comparable to that of our simulation box which represent independent realizations. For each separation bin of width $\Delta s = 0.5 \, h^{-1} \text{Mpc}$ we perform Monte-Carlo realizations of $F(s)$, where $F(s)$ is averaged over the halo/spectrum pairs lying in the given bin. The cosmic variance error is then the $1\sigma$ scatter around $F(s)$. The error of $F(s)$ depends on the total number of galaxies/haloes with separation $s$. A sample of about 40 LOS through our simulation box gives a cumulative halo function $N(<s)$ close to the cumulative galaxy function $N_{\text{LBG}}$ of the ASSP sample for separation $s \lesssim 1 \, h^{-1} \text{Mpc}$. Note also that the number of haloes with $M > M_{\text{min}}$ in the simulation is only about 110. This means that we severely oversample the halo catalogue when we estimate the cosmic variance on large scales. However, at the small scales in which we are mainly interested this should have little effect.

5 THE HALO–FLUX CORRELATION OF OBSERVED AND SIMULATED SPECTRA

5.1 Wind bubbles and the halo–flux correlation

We first explore the shape of $F$ assuming that all wind bubbles have the same radius, $r_w$, independent of the halo mass that they surround. In the right panel of Fig. 3 we plot $F$ for haloes with mass $M > M_{\text{min}}$. We show curves for $r_w = 0.5$, 1 and $1.5 \, h^{-1} \text{Mpc}$ as indicated in the figure. The corresponding filling factors are $f_{\text{H}_\alpha} \sim 0.2$ per cent, $\sim 1.6$ per cent and $\sim 5.3$ per cent respectively, small enough to ensure that the flux PDF, the power spectrum and the line distribution of the Ly$\alpha$ forest are not affected. The dashed-dotted curve corresponds to $F$ in the absence of outflows. The filled symbols show the observed correlation as given in ASSP. The shaded area shows the cosmic variance error, and is calculated (cf. Section 4.3) from the cumulative halo distribution $N(<s)$ also shown as a histogram in the right panel of Fig. 3 (right axis). The cosmic variance error is only shown for $r_w = 0.5 \, h^{-1} \text{Mpc}$.

The trend of decreasing mean flux with decreasing distance to the next galaxy in the observations is well reproduced by the numerical simulations. At small separation this trend is inverted if the wind bubbles are included. Like Kollmeier et al. (2003a,b), we find that the strong increase of the mean flux level at $s \gtrsim 0.5 \, h^{-1} \text{Mpc}$ is only reproduced for bubbles with size as large as $r_w \gtrsim 1.5 \, h^{-1} \text{Mpc}$. We thus confirm the difficulty of other authors in reproducing the mean flux at small separation unless rather large bubbles are included. Bubbles of this size are, however, difficult to explain even with very efficient galactic superwinds (Croft et al. 2002).

There are thus two possibilities: the simple model of spherical wind bubbles of constant size around massive haloes does not describe the effect of winds of observed galaxies properly; or the observed strong increase of the flux level in the ASSP is a statistical fluke. We will explore both possibilities in the following sections in more detail.

5.2 The effect of peculiar motions and thermal broadening

Kollmeier et al. (2003a,b) have pointed out that the peculiar velocities of matter infalling onto the galaxy/DM halo expected from hierarchical structure formation should have a strong influence on $F(s)$ at small separations. In Fig. 4 we investigate the effect of coherent small-scale peculiar velocities and thermal velocities. In the left and middle panel we show $F(s)$ with and without peculiar and thermal velocities for wind bubbles of size $r_w = 0.5 \, h^{-1} \text{Mpc}$ and $r_w = 1 \, h^{-1} \text{Mpc}$, respectively. Without peculiar velocity the increase of the mean flux at small separation is much more pronounced. This is easy to understand. In redshift space the velocity shear of the infalling surrounding material fills in the cavity which is caused by...
the fully ionized wind bubble in real space. For the smaller bubble size such an effect is also visible for the thermal velocities alone. Neutral hydrogen HI lying at the boundary of the ionized bubbles leads to significant Lyα absorption at the redshift halo position when the ratio $b/r_w$ is large enough. The peculiar velocity will, however, always dominate over the effect of thermal velocities. This can be seen in the right panel where we have scaled the temperature of the gas up and down. This has very little effect on $F(s)$ calculated with peculiar velocities even for the smaller bubble size. We have also computed $F(s)$ for various values of the density threshold in equation (6). We found only small differences for $10^{-3} \lesssim b/\rho \lesssim 100$. Such overdensities occur typically within $\lesssim 0.2$ $h^{-1}$ Mpc (comoving) from the halo (Kollmeier et al. 2003a) which is smaller than the bubble radii considered here. Note that the good agreement of the flux PDF of the synthetic spectra with the McDonald et al. (2000) data found in Section 3 disappeared if peculiar velocities were set to zero. This indicates that our simulations have similar peculiar velocities as the gas responsible for the observed absorption – at least in a volume-averaged sense.

In the left panel of Fig. 5 the black dots show the velocity dispersion for haloes of mass $M$. For each DM halo the velocity shear $\sigma_v$ along the line of sight was calculated along three perpendicular LOS directions for an interval $\Delta \pi = 2$ $h^{-1}$ Mpc centred on the halo position. The histogram shows the mean velocity dispersion averaged over bins with $\Delta \log M = 0.25$. The velocity shear clearly increases with halo mass. This trend is present for a range of $\Delta \pi$, $1$ $h^{-1}$ Mpc $\lesssim \Delta \pi \lesssim 2$ $h^{-1}$ Mpc. The middle and right panels of Fig. 5 show the transmitted flux at separation $s < 0.5$ $h^{-1}$ Mpc, $F_0 = F(s < 0.5)$, as a function of the halo mass $M$ assuming that all haloes with $M \gtrsim 10^{11}$ $M_\odot / h$ are surrounded with bubbles of radius $r_w = 0.5$ and $r_w = 1$ $h^{-1}$ Mpc, respectively. As expected $F_0$ is significantly larger for $r_w = 1$ $h^{-1}$ Mpc. The dependence on halo mass appears to be stronger for $r_w = 1$ $h^{-1}$ Mpc than for $r_w = 0.5$ $h^{-1}$ Mpc. In the former case, $F_0$ decreases with increasing $M$ by about $\Delta F_0 = 0.20$ from $M = 10^{11}$ to $M = 10^{12}$ $M_\odot / h$. This is consistent with the trend in the left panel which predicts that low-mass haloes satisfy the condition $r_w \gtrsim (1+z)\sigma_v / H(z)$ more often than massive haloes as a result of the $\sigma_v - M$ dependence. For smaller bubbles, $r_w \lesssim 0.5$ $h^{-1}$ Mpc, this appears to be rarely the case, independent of the mass of the halo.

To investigate the effect of the velocity field near haloes further we show the real-space density and velocity field of the gas, as well as the absorption spectrum along two lines of sight passing haloes with $M = 2 \times 10^{11}$ (left panel) and $3 \times 10^{12}$ $M_\odot / h$ (right panel) in Fig. 6. The solid and dotted lines are for the case with and without wind bubbles, respectively. The radius of the bubbles is $r_w = 0.5$ $h^{-1}$ Mpc for both haloes. The difference between the two cases is striking. For the smaller-mass halo in the left panel the flux level at the position of the halo is $\tilde{F} \sim 0.9$ if the bubble is included, whereas $\tilde{F}$ saturates at the redshift space position of the more massive halo in the right panel. Since the density fields near the haloes in both panels show similar features, the difference in

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### Figure 4

Left and middle panels: the transmitted flux in redshift space (solid), with gas and halo velocities set to zero but with non-zero Doppler parameter $b$ (long-dashed), and with velocities and $b$ set to zero (short-dashed). Panel to the right: the transmitted flux as a function of the mean IGM temperature $\langle T \rangle$ as indicated on the figure. In each panel the horizontal dotted line shows the mean flux level, $(F) = 0.67$.

### Figure 5

Left panel: the velocity dispersion along the line of sight $\sigma_v$ near haloes of mass $M \gtrsim 10^{11}$ $M_\odot / h$. The solid line is $\sigma_v$ averaged over bins of size $\Delta \log M = 0.25$. Middle and right panels: the transmitted flux at separation $s \leq 0.5$ $h^{-1}$ Mpc, $F_0$, as a function of $M$. The horizontal dashed line is $(F) = 0.67$.
Figure 6. One-dimensional gas density, velocity (real space) and transmitted flux (redshift space) along two lines of sight extracted from simulation at $z = 3$, running through the centres of haloes with mass $M = 2 \times 10^{11}$ (left panel) and $M = 3 \times 10^{12} M_\odot / h$ (right panel). The halo positions along the line of sight are shown by vertical lines. The transmitted flux is computed with winds (solid line) and without winds (dotted line). The bubble radius is $r_w = 0.5 \, h^{-1} \text{Mpc}$.

Figure 7. Left panel: the transmitted flux $\bar{F}$ for bubbles of size $r_w = 0.5 \, h^{-1} \text{Mpc}$ after adding a Gaussian deviate of variance $\sigma_z = 0, 150$ and $300 \, \text{km s}^{-1}$ to the LOS halo position. Middle panel: the cumulative halo function (histogram) and the $1\sigma$ error (shaded area) for the model with $\sigma_z = 150 \, \text{km s}^{-1}$. Right panel: the transmitted flux with a mean flux level $\langle F \rangle = 0.62, 0.67$ and $0.72$. The filled symbols are the measurements by Adelberger et al.

5.3 The role of galaxy redshift error, LOS smoothing, and flux normalization

There is a number of other uncertainties which have to be taken into account when comparing the observed $\bar{F}(s)$ with that calculated from the numerical simulation. Most important are the uncertain redshifts of the LBGs. There are large systematic shifts of several hundred km s$^{-1}$ between the redshift of nebular emission lines, Ly$\alpha$ emission and interstellar absorption lines which make the assignment of a redshift somewhat ambiguous. ASSP give $\Delta z \geq 0.002$ or about 150 km s$^{-1}$ as a typical error. We know, however, relatively little about the dynamical state of the responsible gas and there may well be larger systematic errors. The left panel of Fig. 7 shows the effect of adding a Gaussian distributed error to the redshift of the DM haloes. The solid curve shows the same model as the solid curve in the right panel of Fig. 3 (‘massive’ haloes, $r_w = 0.5 \, h^{-1} \text{Mpc}$). Long-dashed and short-dashed curves are for redshift errors of $\sigma_z = 150 \, \text{km s}^{-1}$ and $300 \, \text{km s}^{-1}$, respectively. As expected, introducing redshift errors smoothes out the large depression of the flux level at distances of up to $\sim 2 \, h^{-1} \text{Mpc}$ to the line of sight. This substantially improves the agreement with the ASSP measurements on these scales.

In the middle panel of Fig. 7 the shaded area shows the $1\sigma$ error of the simulated $\bar{F}(s)$ for our ensemble of LBG-size haloes with mass $M \geq 5 \times 10^{11} M_\odot / h$. The error bars show the errors quoted by ASSP. To make the comparison easier the errors are plotted for the model with $\sigma_z = 150 \, \text{km s}^{-1}$ and $300 \, \text{km s}^{-1}$, respectively. The solid histogram shows the cumulative halo function $N(<s)$ (right axis) as a function of separation. At $s \leq 1 \, h^{-1} \text{Mpc}$ the cumulative halo function $N(<s)$ is close to the cumulative galaxy function $N_{\text{LBG}}(<s)$ of the ASSP sample, with $N(s < 0.5) = 3$ and $N(s < 1) = 11$ (respectively, 3 and 12 in the ASSP sample). Our errors are generally larger than that of ASSP, at $s < 0.5 \, h^{-1} \text{Mpc}$ by about 30 per cent. At larger scales
5.4 The effect of varying the bubble model

5.4.1 \( F \) with halo-mass-dependent bubble size

So far we have considered bubbles with sizes that are independent of mass. Unfortunately, it is far from obvious what scaling with mass is appropriate. Obviously, on average more stars and a larger total energy input from stellar feedback are expected in the more massive haloes. However, simple binding energy arguments show that winds should escape more easily from low-mass haloes (Larson 1974; Dekel & Silk 1986). This may be properly taken into account that winds should escape more easily from low-mass haloes (Larson 1974; Dekel & Silk 1986). This may be properly taken into account.

The left panels of Fig. 8 compare the case of constant bubble size with those assuming a power-law relation \( r_w(M) \propto M^{v} \) with \( v = +1/2 \) and \(-1/2\), at the minimal halo mass \( M_{\text{min}} = 5 \times 10^{11} \, M_{\odot}/h \). The top panels are for \( r_w = 0.5 \) and the bottom panels are for \( r_w = 1.0 \). For \( v = +1/2 \), \( F(s) \) is very sensitive to the value of \( r_w \) and the mean flux at \( s < 1 \, h^{-1} \text{Mpc} \) rises to values as large as \( F = 0.8 \) for the larger of the two bubble sizes, improving the agreement with the results of ASSP. However, this is at the cost of the good agreement with the data on larger scales. It thus appears difficult to reproduce the observations of ASSP at small and large scales simultaneously by changing the mass–radius relation of the bubbles.

In the right panels of Fig. 8 we investigate the effect of changing the minimum mass above which DM haloes are assumed to be surrounded by wind bubbles. Changing \( M_{\text{min}} \) will change the total space density of the halo as given in Table 2. There is less absorption if \( M_{\text{min}} \) decreases, mainly as a result of the increasing bubble filling factor. With a bubble radius \( r_w = 1 \, h^{-1} \text{Mpc} \), the filling factor is \( f_{\text{fill}} \sim 15 \) per cent and \sim 1 per cent for \( M \geq 10^{13} \) and \( M \geq 10^{12} \, M_{\odot}/h \), respectively.

5.4.2 The starburst model

So far we have assumed that LBGs are hosted by the most massive haloes available and have introduced a minimum mass to match their space density with that observed for LBGs. This assumes that LBGs are long-lived and have the same brightness over a fair fraction of the Hubble time. This picture may be in conflict with the lack of the evolution of the space density of bright LBGs towards very high redshift which is more easily explained if LBGs are starbursts with a duty cycle of high star-formation rate of less than 10 per cent. This would mean that there is an about a factor 10 or more larger number of DM haloes needed to host these LBGs than in the massive halo picture and the DM haloes hosting LBGs are less massive on average. As should be apparent from the previous section, for a given bubble size surrounding a typical LBG halo, this should reduce the effect of the velocity shear. It is then easier to reproduce the high mean flux levels at small separation. We investigate this alternative ‘starburst’ picture here in more detail. Kolatt et al. (1999) developed a simple, ad hoc procedure to assign LBGs to colliding haloes identified in their N-body simulation and were able to reproduce both the number density and clustering properties of LBGs. Here, we will do something simpler and randomly select haloes with mass \( M \geq 10^{11} \, M_{\odot}/h \) such that their number density matches the observed LBGs number density, \( n_{\text{LBG}} \sim 0.004 \, h^2 \text{Mpc}^{-3} \). This corresponds to a duty cycle of 10 per cent. Note that only 10 per cent of the haloes with a mass above \( 10^{11} \, M_{\odot}/h \) are haloes as massive as our ‘massive halo’ picture, i.e. have masses \( M \geq 10^{12} \, M_{\odot}/h \) (cf. Table 2). Most of the LBG-size haloes have \( M \lesssim 5 \times 10^{11} \, M_{\odot}/h \) in the ‘starburst’ scenario. We have produced 120 different Monte Carlo realizations of such catalogues of starbursting haloes. In the left panel of Fig. 9 we compare \( F(s) \) for the ‘massive halo’ model (solid curve) and the ‘starburst’ picture (dashed curve). For both we assume a bubble radius \( r_w = 1 \, h^{-1} \text{Mpc} \) and a redshift error of the halo position of \( \sigma_z = 150 \, \text{km} \, \text{s}^{-1} \). At \( s \lesssim 1 \, h^{-1} \text{Mpc} \), \( F(s) \) is larger in the ‘starburst’ model by about 20 per cent. We also plot the 1σ error for the starburst model and the ASSP measurements. Our ‘starburst’ model with \( r_w = 1 \, h^{-1} \text{Mpc} \) and a volume filling factor of 2 per cent appears to be fully consistent with the ASSP measurement. Even at \( s \lesssim 0.5 \, h^{-1} \text{Mpc} \), where our estimate of the cosmic variance error is 30 per cent larger than that of ASSP, \( F(s) \) in the starburst model is only slightly more than 1σ below the ASSP measurement despite a difference in the mean flux of \( \Delta F(s) = 0.3 \). Note that for the quoted volume filling factor we only took into account the bubbles around the haloes that host LBGs at a given time. If the bubbles last longer than the LBG phase the volume filling factor would be higher by the same factor.
5.4.3 The clustering of LBG haloes

As discussed above our procedure to identify the host haloes of LBGs in the ‘starburst’ model is rather crude. Unlike Kolatt et al. (1999) and Wechsler et al. (2001) we do not identify ‘colliding’ haloes but choose a fraction of less massive haloes at random. It is thus not obvious that the host haloes chosen in this way will reproduce the clustering properties of LBGs. In the right panel of Fig. 9 we compare the two-point correlation function of our ‘starburst’ and ‘massive halo’ models. The correlation function in the ‘starburst’ model was computed from 120 different Monte Carlo realizations of catalogues of starbursting haloes. The correlation lengths for the ‘starburst’ and ‘massive halo’ models are 2.5 h⁻¹ Mpc, and 3.5 h⁻¹ Mpc, respectively. The observed correlation length is \( r_0 \approx 3-4 \) h⁻¹ Mpc (Adelberger et al. 1998; Arnouts et al. 1999; Adelberger 2000; Gavilánisco & Dickinson 2001). The massive halo picture is thus not obvious that the host haloes chosen in this way will reproduce the observed clustering properties if they select ‘colliding haloes’ as host galaxies of starbursting LBGs. We thus do not investigate this issue any further here.

5.5 Cosmic variance

To understand better why the cosmic variance error becomes so large at small separation we plot the probability distribution \( P(F_0) \) in Fig. 10. The distribution is very broad. About 40 per cent of the haloes have \( F_0 \) larger than 0.8. The probability of finding three (uncorrelated) haloes with \( F \geq 0.8 \) at separation \( s \leq 0.5 \) h⁻¹ Mpc is thus still \( \sim 5-10 \) per cent. The symbols at the bottom of Fig. 5 show the increase of \( F_0 \) if the thermal motions and the velocity shear are artificially reduced by a factor of two. The broad probability distribution also explains why the suggestion of Croft et. (2002) that LBGs are preferentially located in haloes with a low-density environment has such a strong effect on the mean transmission.

6 DISCUSSION

We used mock spectra obtained from numerical simulations to investigate the effect of fully ionized bubbles around the galactic haloes hosting LBGs on the \( \text{H}_\text{i} \) distribution in the IGM as probed by the Ly\( \alpha \) forest in QSO spectra. By matching the probability distribution and the power spectrum of the flux of mock spectra obtained from an \( N \)-body simulation of a \( \Lambda \)CDM model with those of observed spectra we derive an upper limit of 10 per cent on the volume filling factor of bubbles. We further considered the mean transmitted flux...
at a given distance from galactic haloes. To model the observational data we computed the mean transmitted flux $\bar{F}$ from an annular average of the two-dimensional conditional distribution $F(\sigma_x, \sigma_y)$ in the same way as done by Adelberger et al. (2003). As in previous investigations we find that the observed decrease of the transmission at separation $1 \leq s \leq 5 h^{-1}$ Mpc to a halo is well reproduced by the simulated spectra. It is due to the increased matter density around haloes. We further find that an increase of the transmissivity at the smallest separations $s \leq 0.5 h^{-1}$ Mpc requires bubble radii of $r_w = 1.5 h^{-1}$ Mpc. This is two to three times larger than the separation at which the increase of the transmissivity is observed. As discussed by Kollmeier et al. (2003a,b) this is because in redshift space the velocity shear of the material infalling on to the halo fills in the cavity which is caused by the fully ionized wind bubble in real space. For a model where the haloes are surrounded with bubbles of size $r_w = 1.5 h^{-1}$ Mpc there is no effect on either the flux PDF or the flux power spectrum. The flux PDF and the flux power spectrum of our simulated spectra are fully consistent with those of observed spectra. This is because the volume filling factor of the ionized bubbles is still small. We have investigated a variety of scalings of bubble size with halo mass. We also compared a model where LBGs are long-lived and are hosted by the most massive haloes with a model where LBGs are short-lived starbursts and are hosted by more numerous less massive haloes. Unlike Kolatt et al. (1999) and Wechsler et al. (2001), we do not identify starbursts as ‘colliding haloes’ but choose a fraction of less massive haloes at random. Matching the observed transmission at scales $1 \leq s \leq 5 h^{-1}$ Mpc is not difficult and there is certainly more than one way to do so. Introducing redshift errors at the expected level improves the agreement with the observations. The increase in transmission at small separation in our model spectra is, however, generally more gradual than that of the observed measurements. Like other authors, we found it challenging to match the observations at the smallest separation. Numerical simulations including galactic winds and simple analytical estimates such as the one in Section A3 in Appendix A agree that the required bubble sizes are difficult to achieve (Bruscoli et al. 2003; Croft et al. 2002). We have thus particularly looked into models which are consistent with the observed measurements and are most ‘economical’ in terms of bubble size. The velocity shear around haloes increases with increasing mass. Haloes of lower mass thus require smaller bubble radii to explain the observed increase of the transmission at $s \leq 0.5 h^{-1}$ Mpc. This may favour a model where LBGs are starburst galaxies hosted by more numerous less massive haloes. We have taken special care in estimating the errors on the flux near galaxies in a similar way as done in the observations. For small separations we found that the expected error due to cosmic variance is about 30 per cent larger than estimated by ASSP. Our starburst model is consistent with all measurements for a bubble radius of $r_w = 1 h^{-1}$ Mpc and a volume filling factor of 2 per cent. However, a bubble size of $r_w = 1 h^{-1}$ Mpc still appears difficult to achieve for galactic winds. If for some reason the velocity shear around haloes is smaller than suggested by our (DM only) simulation this would further reduce the required bubble size and filling factor. The correlation length for our ‘starburst’ model is $2.5 h^{-1}$ Mpc, slightly lower than the observations. Identifying the LBGs with satellites infalling on high-mass haloes, as in Kolatt et al. (1999), would further improve the agreement with the observed clustering properties of LBGs. However, this would be at the expense of matching the observed small-scale flux $\bar{F}$, since the velocity shear increases with halo mass. Therefore we advocate the model in which starbursts occur in small-mass haloes. In a $\Lambda$CDM cosmology of normalization $\sigma_8 = 0.9$, this model yields a reasonable match to the observed clustering of LBGs and transmitted flux $\bar{F}$. The mean transmitted flux of our simulated spectra at the smallest separation is 30 per cent smaller than that of the ASSP sample. We would thus expect the observed mean transmission to decrease substantially for larger samples if the model is correct. A sample with 12 LBGs at separation $s \leq 0.5 h^{-1}$ Mpc should reduce the error to $\sigma \bar{F} \sim 0.1$. This should help to discriminate against alternative models like the suggestion of Croft et al. (2002) that LBGs are located preferentially in low-density environments.

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APPENDIX A: EVOLUTION OF BUBBLES

To explain the lack of H I close to the haloes, the dilute, highly ionized bubble has to survive for a sufficiently long time. In this appendix we present general order-of-magnitude estimates of the wind/shock energetics, relate that to the observed speeds, and estimate cooling time-scales. We heavily rely on the work of Sedov (1959), Weaver et al. (1977), Ostriker & McKee (1988) and Tegmark et al. (1993).

A1 Wind energetics

The main characteristics of the wind, such as the temperature of the inner plasma, the size of the bubble swept up by the shock and the wind velocity, depend on the total energy $E_w$ released in the outflow. Neglecting energy losses, $E_w$ can be estimated as

$$E_w = \frac{1}{2} M_w v_w^2$$

$$\sim 4.3 \times 10^{56} \text{erg} \left( \frac{r_w}{0.1 \text{ h}^{-1} \text{Mpc}} \right)^3 \left( \frac{v_w}{100 \text{ km s}^{-1}} \right)^2 \Delta_g, \quad (A1)$$

where $M_w$ is the gas mass inside $r_w$. To get the numerical estimation we have assumed that $M_w = (4\pi/3) r_w^3 \rho_b \Delta_g$, where $\rho_b$ is the mean background density and $\Delta_g$ is the density contrast inside $r_w$. The characteristic radius $r_w$ and velocity $v_w$ of the wind are poorly constrained, but recent observations suggest that $r_w \gtrsim 0.1 \text{ h}^{-1} \text{Mpc}$ and $v_w \gtrsim 100 \text{ km s}^{-1}$ (Heckman et al. 2000; Pettini et al. 2001; Adelberger et al. 2003). On scales $r \sim 0.1 \text{ h}^{-1} \text{Mpc}$ the value of $\Delta_g$ is typically about 2–5, which gives $E_w \sim 10^{57} \text{erg}$. This value is consistent with the amount of energy $E_{SN}$ released by supernovae during a burst of duration $t_{burst} \sim 10^7 \text{yr}$ and of star-formation rate $M_* \sim 10 M_\odot \text{yr}^{-1}$.

To assess whether the gas will be expelled from the halo one can apply a simple binding argument (e.g. Dekel & Silk 1986). On the one hand, the potential energy $U$ of the gas which lies within the halo is $U \propto M_{halo} v^2 / r_{M}$, where $r_{M}$ is the characteristic radius of the halo, $r_{M} \propto M_{halo}^{1/3}$, and $M_{halo}$ is the mass of the gas which lie within a radius $r_{M}$. The mass of the gas in the halo is $M_{halo} \propto f_g M$, where $f_g$ is the star-formation efficiency and $f_b$ the average fraction of baryons in the halo. The potential energy is then $U \propto M_{halo}^{5/3}$. On the other hand, if one assumes that winds are driven by supernovae, the energy released in the wind is $E_w \propto M_{halo} v^4$, which are then the star-formation rate. Hence, according to this crude estimate, the ratio of the kinetic to binding energy is $E_w / U \sim M_{halo}^{2/3}$. We therefore expect $r_w$ to be a function of $M_{halo}$, and outflows to escape more easily from low-mass haloes (e.g. from dwarf galaxies). On this latter point, the measurements of Adelberger et al. (2003) seem to indicate that, if outflows are the cause of lack of H I absorption in the observed transmitted flux $F$, they can escape from high-mass galaxies as well, and affect the IGM properties on comoving scale as large as 0.5 $h^{-1}$ Mpc.

A2 Propagation of an adiabatic shock

A wind injects energy in the IGM, driving a shock which heats up and collisionally ionizes the material it encounters. We briefly review now the evolution of the shock front and the ionized bubble it encompasses, before radiative losses become important. We assume that the gas follows the perfect gas law $P = \rho T$ where $P$, $\rho$ and $T$ are, respectively, the pressure, the internal energy and the density of the gas, and $\gamma$ is the ratio of specific heats, which is $\gamma = 5/3$ for a monatomic gas. If the expansion of the Universe can be ignored, the shock is well described by a strong adiabatic blast wave which propagates into the IGM. In this regime, accurate self-similar solutions exist. Let $r_{sh}(t)$ be the position of the shock front at time $t$ after the burst. The shock conditions relate the density $\rho_{IGM}$ of the unperturbed IGM to the density $\rho_{sh}$, the pressure $P_{sh}$ and the temperature $T_{sh}$ at the shock front. For a strong adiabatic shock the jump conditions for pressure and density can be expressed as

$$\rho_{sh} = \frac{(\gamma + 1)}{(\gamma - 1)} \rho_{IGM}, \quad P_{sh} = \frac{2}{\gamma - 1} \rho_{IGM} \frac{dr_{sh}}{dt}. \quad (A2)$$

Hence, $\rho_{sh} = 4 \rho_{IGM}$ for an adiabatic index $\gamma = 5/3$. Assuming that the flow is self-similar, $r_{sh}$ can be expressed as a function of the similarity variable $\xi$. The position of the shock front corresponds then to a fixed value of $\xi \simeq 1$. Hence, $r_{sh}$ is given by (Sedov 1959)

$$r_{sh} = \left( \frac{E_w}{\rho_{IGM}} \right)^{1/5} \xi^{2/5}, \quad (A3)$$

where $E_w \sim 10^{57} \text{erg}$ is the energy released in the wind. Inserting the value of $E_w$ in the expressions of $r_{sh}$ and $T_{sh}$, we find

$$r_{sh}(t) \simeq 19 \text{ h}^{-1} \text{kpc} \left( \frac{t}{10^7 \text{yr}} \right)^{0.4} \times \frac{1}{\Delta_{0.2}^2} \left( \Omega_b h^2 \right)^{-0.2} \left( \frac{E_w}{10^{57} \text{erg}} \right)^{0.2} \left( \frac{1 + z}{4} \right)^{-0.6}. \quad (A4)$$

$$T_{sh}(t) \simeq 2.7 \times 10^6 \text{K} \left( \frac{t}{10^7 \text{yr}} \right)^{-1.2} \times \frac{1}{\Delta_{0.2}^2} \left( \Omega_b h^2 \right)^{-0.4} \left( \frac{E_w}{10^{57} \text{erg}} \right)^{0.4} \left( \frac{1 + z}{4} \right)^{-1.2}. \quad (A5)$$

Since the temperature behind the shock front $T \propto P / \rho$ is a decreasing function of $r$, $T$ reaches its minimum at the shock front. This
solution holds only in the regime where the expansion is negligible, i.e. $t \ll t_H$ where $t_H$ is the Hubble time-scale. Self-similar solutions for $t > t_H$ including this latter complication exist (Ostriker & McKee 1988), but since energy and momentum conservation break down for $t \sim t_H$, it is difficult to express the asymptotic solution as a function of the initial $E_w$.

### A3 Cooling time-scales and survival of bubbles

The evolution of such an expanding shock depends critically on cooling which will lower the temperature of the shell and its interior, thereby increasing the amount of H$_1$ hydrogen available for Ly$\alpha$ resonant scattering. The expanding shell cools due to radiative losses. The shock will undergo a phase of shell formation which will occur when the energy dissipated in a volume $\propto r_{sh}(t)^3$ will be of the order of $E_w$.

Cooling is expected to be more efficient at the shock front since, at $T = r_{sh}$, the density (temperature) is larger (lower) than at $r < r_{sh}$. Thus, shell formation first occurs at the shock front. As long as we are in the regime of low density, $\rho \ll \rho_{\text{HII}}$, and high temperature, $T \gg 10^6$ K, radiative cooling is dominated by thermal bremsstrahlung. However, for temperature $T \lesssim 10^6$ K, atomic line cooling might contribute significantly to the cooling rate $\Lambda$ (in erg cm$^{-3}$ s$^{-1}$) if the wind is enriched with heavy elements. For a wind of metallicity $Z \sim 0.1 Z_\odot$ (e.g. Adelberger et al. 2003), the cooling rate is $\Lambda \sim 10^{-22}$ erg cm$^{-3}$ s$^{-1}$ in the range $10^5 \lesssim T \lesssim 10^6$ K (e.g. Sutherland & Dopita 1993).

We assume that cooling at the shock front, where the temperature is $T \lesssim 10^6$ K, is dominated by atomic line transitions. This is true when the temperature of the shell is $T_{sh} \lesssim 10^6$ K, i.e. for times $t \gtrsim 10^6$ yr. The cooling time-scale at the shock front is then $t_{cool}(r_{sh}, t) = 3k_B T_{sh}/(2\Lambda_{22} \Lambda)$,

$$t_{cool}(r_{sh}, t) = 1.9 \times 10^{10} \text{yr} \left(\frac{t}{10^6 \text{yr}}\right)^{-1.2} \Lambda_{22}^{-1} \times \frac{1}{\Delta_g^4} \left(\frac{\Omega_0 h^2}{0.019}\right)^{-1.4} \left(\frac{E_w}{10^{37} \text{erg}}\right)^{0.4} \left(\frac{1 + z}{4}\right)^{-4.2},$$

where the cooling rate $\Lambda_{22}$ is in units of $10^{-22}$ erg cm$^{-3}$ s$^{-1}$ and is assumed to be constant in the temperature range $10^4 \lesssim T_{sh} \lesssim 10^6$ K.

Once the shell starts forming, the temperature $T_{sh}$ of the shell cools to $\sim 10^5$ K on a very short time-scale. However, in the range $T \lesssim 10^4$ K, the cooling rate $\Lambda$ drops sharply from $\sim 10^{-22}$ down to $\sim 10^{-26}$ erg cm$^{-3}$ s$^{-1}$ (Sutherland & Dopita 1993) and the shell cannot cool below temperature $T \sim 10^3$ K. We now define a radius $r_{cool}(t)$ such that the cooling time-scale $t_{cool}$ as a function of radius $r$ satisfies $t_{cool} < t$ for $r > r_{cool}$. According to this definition, the gas between $r_{cool}(t)$ and $r_{sh}(t)$ has cooled on to the shell by the time $t$. To express $t_{cool}$ as a function of $r$ and $t$, it should be noted that the number density $n_w(r, t)$ behind the shock front (in cm$^{-3}$) scales as $n_w(r, t) = (r/r_{sh})^n n_w(r, t)$ with $n \sim 4.5$ typically. This relation is, however, valid in the asymptotic regime only. Therefore, assuming that the plasma has reached pressure equilibrium, the temperature $T_{sh}(r, t)$ behind the shock front is $T_{w} = (r_{sh}/r)^\alpha T_{sh}$, and the radiative cooling time-scale $t_{cool}$ is

$$t_{cool}(r, t) = \left(\frac{r_{sh}}{r}\right)^{2n} t_{cool}(r_{sh}, t).$$

In Fig. A1 the solid curves are $r_{cool}(t)$ for various cooling time rates $\Lambda_{22}$, and the dashed curve is $r_{sh}(t)$. The shell starts forming at the shock front when $r_{cool}(t)$ intersects $r_{sh}(t)$. For larger time, cooling occurs for $r_{cool}(t) < r < r_{sh}(t)$ (shaded area). Since most of the gas swept up by the wind lies at $r \lesssim r_{sh}$, a large fraction of gas cools down to $T \sim 10^4$ K shortly after the shell formation. For $\Lambda_{22} = 1$, about 30 per cent of the mass enclosed in the bubble has cooled into the shell over a Hubble time $t_H$, where $t_H \approx 1.2 \times 10^7 (\Omega_0 h^2)^{-1/2}$ at $z = 3$. For a high-metallicity wind $Z \gtrsim 0.1 Z_\odot$, cooling becomes severe for $t \gtrsim 10^7$ yr, a time-scale shorter than the Hubble-time scale $t_H$ at $z = 3$. On the contrary, for a low-metallicity wind, $Z \ll 0.1 Z_\odot$, the cooling rate is $\Lambda_{22} \ll 1$ and cooling will not become significant for $t \lesssim t_H$. Hence, most of the bubbles will not survive over a Hubble time-scale unless the cooling time rate $\Lambda_{22}$, or the wind metallicity $Z$, is very low. For radii $r \lesssim r_{cool}(t)$, the temperature is higher than $10^5$ K, and thermal bremsstrahlung and inverse Compton cooling prevail over line cooling, but the corresponding cooling times are still smaller than the Hubble time at $z = 3$.

### A4 Absorption by the cold shell

To compute the thickness $\Delta s(t)$ of the shell, we assume that the pressure is constant through the inner plasma and the shell. Since the temperature of the shell, $T_{sh} \sim 10^5$ K, is much lower than the overall temperature $T_w > 10^6$ K of the highly ionized plasma, we expect $\Delta s(t) \ll r_{sh}(t)$, i.e. the shell is thin. Taking into account this assumption, we find

$$\Delta s = \frac{1}{3} \left(\frac{T_{sh}}{T_w}\right) \left[1 - \left(\frac{r_{cool}}{r_{sh}}\right)^3\right],$$

a result consistent with a thin-shell approximation. Assuming $T_w = 100 T_{sh} = 10^6$ K and neglecting the expansion of the Universe gives $\Delta s \lesssim 0.01 r_{sh}$. $\Delta s$ should thus not exceed a few $h^{-1}$ kpc, a scale which cannot be resolved in our simulation. The optical depth for resonant Ly$\alpha$ scattering is given by

$$\tau \sim \sigma_{\text{Ly}\alpha} \Delta s \sim 1.6 \pi \Delta_g \left(\frac{\Omega_0 h^2}{0.019}\right) \left(\frac{1 + z}{4}\right) \left(\frac{r_{sh}}{0.1 \ h^{-1} \text{Mpc}}\right).$$

Figure A1. The cooling radius $r_{cool}(t)$ (solid curves) for various cooling time rates. The dashed curve is $r_{sh}(t)$. 

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where \( \eta \simeq 1 \) is the HI neutral fraction at temperature \( T \sim 10^4 \) K. Ly\( \alpha \) absorption is expected for reasonable values of the density contrast \( \Delta g \sim 2-3 \), the typical comoving radius of the bubble \( r_w \sim 0.1 \) \( h^{-1} \) Mpc, and the baryon fraction \( \Omega_b h^2 \sim 0.02 \). However, since we cannot resolve the thin shell in our simulation, we have nevertheless ignored its Ly\( \alpha \) absorption.

**APPENDIX B: NUMERICAL CONVERGENCE**

To assess the sensitivity of the transmitted flux to numerical resolution, we perform a calculation with the highest-resolution simulation of the series, which has a mass resolution \( m = 1.66 \times 10^8 \) M\(_{\odot}/h \). The halo catalogue and the flux in the simulation are computed as in Section 3 with a pixel resolution \( \Delta = 60 \ h^{-1} \) kpc (comoving). We added a Gaussian deviate of variance \( \sigma_v = 150 \) km s\(^{-1} \) to the halo positions. In Fig. B1 we show the transmitted flux as computed in the ‘starburst’ scenario for a bubble radius \( r_w = 1 \) \( h^{-1} \) Mpc. The mass resolution is \( m = 9.52 \) (dashed curve) and \( m = 1.66 \times 10^8 \) M\(_{\odot}/h \) (solid curve). In the latter the strong absorption lines are better resolved, thereby slightly reducing the normalization of the optical depth \( A \). Therefore, the optical depth of the absorbing material is lower, and the transmitted flux is slightly larger, by about \( \Delta F \sim 0.02 \) on scales of \( 2 \leq s \leq 4 \) \( h^{-1} \) Mpc.

![Figure B1](https://example.com/fig_b1.png)

**Figure B1.** The transmitted flux \( F \) in the ‘starburst’ scenario with bubbles of radius \( r_w = 1 \) \( h^{-1} \) Mpc, with DM particle mass resolution \( m = 9.52 \times 10^8 \) (dashed curve) and \( 1.66 \times 10^8 \) M\(_{\odot}/h \) (solid curve). The horizontal dashed line is the flux normalization, \( \langle F \rangle = 0.67 \).

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