OXYGEN PUMPING. II. PROBING THE INHOMOGENEOUS METAL ENRICHMENT AT THE EPOCH OF REIONIZATION WITH HIGH-FREQUENCY CMB OBSERVATIONS

Carlos Hernández-Monteagudo, Zoltán Haiman, Licia Verde, and Raul Jimenez

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ABSTRACT

At the epoch of reionization, when the high-redshift intergalactic medium (IGM) is being enriched with metals, the 63.2 μm fine-structure line of O i is pumped by the ~1300 Å soft UV background and introduces a spectral distortion in the cosmic microwave background (CMB). Here we use a toy model for the spatial distribution of neutral oxygen in which metal bubbles surround dark matter halos, and compute the fluctuations of this distortion and the angular power spectrum it imprints on the CMB. We discuss the dependence of the power spectrum on the velocity of the winds polluting the IGM with metals, the minimum mass of the halos producing these winds, and the cosmic epoch when the O i pumping occurs. We find that, although the clustering signal of the CMB distortion is weak [(εy)_{rms} ≤ 10^{-7}; roughly corresponding to a temperature anisotropy of ~1 nK], it may be reachable in deep integrations with high-sensitivity infrared detectors. Even without a detection, these instruments should be able to set useful constraints on the heavy-element enrichment history of the IGM.

Subject headings: atomic processes — cosmic microwave background — cosmology: theory — intergalactic medium

1. INTRODUCTION

Probing the Dark Ages—the epoch between the last scattering surface of the cosmic microwave background (CMB) at z ~ 1000 and the completion of the reionization of the intergalactic medium (IGM) at z ~ 6, including the formation of the first luminous objects—constitutes the next frontier of observational cosmology. Among the major open questions are the nature of the objects that reionized the universe and the origin of the first heavy elements, as well as the efficiency with which they were mixed into the high-redshift IGM.

To date, the reionization history has been constrained by observations of Lyman-series absorption spectra of z ~ 6 quasars (see Fan et al. 2006 for a recent review), by gamma-ray burst observations at high redshifts (Totani et al. 2006), by observations of Lyα-selected galaxies (Kashikawa et al. 2006; McQuinn et al. 2007; Dijkstra et al. 2007; Malhotra & Rhoads 2006), and by CMB polarization anisotropies (Kogut et al. 2003; Page et al. 2007; Spergel et al. 2007). The present constraints, however, are consistent with a wide range of scenarios, and considerable theoretical and experimental effort has been devoted to developing new probes of the high-redshift universe. A promising possibility is to use the redshifted 21 cm hyperfine line of H i (see Furlanetto et al. 2006 and references therein). Metal enrichment of the high-redshift IGM may be another useful tracer of the reionization process. For example, Basu et al. (2004) and Hernández-Monteagudo et al. (2006) considered the elastic resonant scattering of CMB photons by intergalactic metals. Recent studies have focused, in particular, on the detectability of neutral oxygen (O i; Oh 2002; Basu et al. 2004), as this element is thought to have been produced in abundance by the first stars (Heger & Woosley 2002). Scattering of UV photons by O i, and the corresponding absorption features in the spectra of quasars—the O i forest—was proposed as a possible observable by Oh (2002) and may have recently been detected at z ~ 6 (Becker et al. 2006).

Since O i and H i are in charge-exchange equilibrium, oxygen is likely to be highly ionized in regions where hydrogen is actively ionized. However, the recombination time for oxygen is shorter than the Hubble time, and it can be neutral even in “fossil” H ii regions where hydrogen has been ionized, but where short-lived ionizing sources have turned off, allowing the region to recombine (Oh 2002). Oh & Haiman (2003) show that the filling factor of such fossil H ii regions can be large (≥50%) prior to reionization.

In a previous paper (Hernández-Monteagudo et al. 2007, hereafter Paper I), we showed that the 63.2 μm fine-structure line of neutral O i can be pumped by the ~1300 Å soft UV background, via the Balmer-α line of O i. This is analogous to the Wouthuysen-field effect (Wouthuysen 1952) for exciting the 21 cm line of cosmic H i. In Paper I we found that O i at redshift z should be seen in emission at (1+z)63.2 μm, and for 7 < z < 10 it would produce a mean spectral distortion of the CMB with a y-parameter of

\[
y = \frac{\Delta I_y}{B_y(T_{CMB})} = (10^{-9} - 3 \times 10^{-8}) \left( \frac{Z}{10^{-3} Z_\odot} \right) \left( \frac{I_{UV}}{10} \right),
\]

where Z is the mean metallicity of the IGM and I_{UV} is the UV background intensity at 1300 Å in units of 10^{-21} erg s^{-1} Hz^{-1} cm^{-2} sr^{-1}. In principle, this distortion could be detectable through a precise future measurement of the CMB spectrum (see Fixsen & Mather [2002] for the prospects of such measurements), and would then open the possibility of performing tomography of the metal distribution. In combination with H i 21 cm studies, it could yield direct measurements of the abundances and spatial distribution of metals in the high-redshift IGM.

Since oxygen pollution at high redshift is likely associated with most overdense regions in the IGM, hosting the first star-forming...
activity, there must be spatial fluctuations in the O I abundance, causing fluctuations in the corresponding $\nu$-distortion. In this paper we consider the clustering properties of this signal. In particular, we use toy models to describe the metal distribution and compute the two-dimensional angular power spectrum of the CMB intensity. This would be appropriate for an instrument with a single (or a few discrete) frequency bands that probes the metals in a single (or a few discrete) narrow redshift bin. We find that the clustering signal induced by the inhomogeneous O I pumping is small, but possibly detectable with forthcoming instruments, e.g., with deep integrations (for 2–3 months) with the full Atacama Large Millimeter Array (ALMA). A detection of the metal enrichment during the Dark Ages would be complementary to $H_\perp$ 21 cm studies, and would provide clues about the distribution of metals in the IGM before galaxy formation started in full.

The rest of this paper is organized as follows. In §2 we discuss the basics of the O I pumping process. In §3 we outline the computation of the distortion induced by O I pumping along a given line of sight in the CMB. In §4 we discuss our method to compute the angular power spectrum of the distortion, with the basic assumption that metals trace the distribution of collapsed halos. In §5 we introduce toy models to describe the spatial distribution of metals, with the basic assumption that metals cluster around dark matter halos. In §6 we present our main result on the angular power spectrum, and discuss its dependence on the basic model parameters, as well as its detectability. Finally, in §7 we summarize the implications of this work and offer our conclusions. Throughout this paper we adopt a set of “concordance” cosmological parameters for a flat universe, $\Omega_m = 0.29$, $\Omega_{\Lambda} = 0.71$, $\Omega_b = 0.047$, and $h = 0.72$, with a power spectrum normalization $\sigma_8 = 0.75$ and slope $n = 0.99$.

2. THE BALMER-\( \alpha \) PUMPING OF O I

In Paper I we computed the distortion in the CMB induced by neutral oxygen in environments where UV radiation is pumping the fine-structure 63.2 $\mu$m M1 transition (between the $n = 2$ electronic states $^3P_2$ and $^3P_1$) via the Balmer-\( \alpha \) line at $\sim 1302$ Å, connecting these two states with the excited $n = 3$ electronic state $^3S_1$. In what follows we denote the $^3P_2$ and $^3P_1$ states as “0” and “1,” respectively, and the excited state $^3S_1$ as “2.” Energy and frequency differences between levels will be denoted as $E_{ij}$ and $\nu_{ij}$, with $i = 0$, 1, 2, and $j > i$.

In Paper I we showed that this process of Balmer-\( \alpha \) pumping modifies the occupation of the levels 0 and 1, and therefore introduces a shift in the spin temperature, $T_S$, so that it slightly departs from the CMB temperature, $T_{\text{CMB}}$. Since at reionization the UV background flux at $\nu_{10}$ is smaller than at $\nu_{21}$, $T_S$ will be above $T_{\text{CMB}}$, producing an excess of 63.2 $\mu$m photons. Each of these excess photons is generated by the sequence of three radiative transitions $0 \rightarrow 2 \rightarrow 1 \rightarrow 0$, i.e., a photon of frequency $\nu_{20}$ is broken into two photons of frequencies $\nu_{21}$ and $\nu_{10}$.

In Paper I we were interested in the average distortion over the full sky, and considered the optically thin limit for the Balmer-\( \alpha \) transition. While this is justified if the O I is homogeneously distributed (in which case $\tau_{20} \sim 0.01 Z/(10^{-2.5} Z_\odot)$ and $\tau_{21} \ll 1$), here we consider metal distributions that can be highly inhomogeneous. The metal-enriched bubbles can potentially contain a large concentration of metals and become optically thick at $\nu_{20}$. For this reason, we consider the case $\tau_{20} \gtrsim 1$ and self-consistently solve the equations of radiative transfer to follow the UV intensity as the background radiation penetrates the O I-rich bubble.

![Fig. 1.—Distortion parameter $\nu$ induced by the Balmer-\( \alpha \) pumping of the O I 63.2 $\mu$m transition as a function of the optical depths $\tau_{10}$ and $\tau_{20}$ at redshift $z = 7$. The optically thin solution presented in Paper I is shown as the solid curve, accurately describing the system at low and intermediate $\tau_{10}$. When the 2 $\rightarrow$ 0 transition becomes optically thick ($\tau_{20} \sim 1$, corresponding to $70 < 3 \times 10^{-4}$), the results are obtained by tracking the photon and level populations as dictated by eq. (2), and as shown by the dashed curve. Note that at low $\tau_{10}$ there is a slight mismatch between the solid and dashed lines, caused by numerical error due to the finite size of our integration step.](image)
3. The Spectral Distortions in the CMB

The total distortion introduced by $O_1$ in the CMB at a given observing frequency $\nu_{\text{obs}}$ is

$$y = \int dr \frac{j_\nu}{B_\nu(T_0)} r \approx \frac{j_\nu}{B_\nu(T_0)} \nu_{\text{obs}} \Delta r,$$  

(3)

where $j_\nu$ is an effective emissivity which is nonzero only at the redshift satisfying $1 + z_i \approx \nu_{\text{obs}} / \nu_{\text{obs}}$ (see Paper I for more details). For a uniform $O_1$ distribution, and assuming that the $O_1$ atoms follow the Hubble expansion, the effective length $\Delta r$ along which this emissivity is not zero is $cH^{-1}(z_i) \Delta \nu_{\text{obs}}/\nu_{\text{obs}}$, where $\Delta \nu_{\text{obs}}$ is the width of the line. However, if oxygen is clumped (for example, confined within bubbles) then the effective length $\Delta r$ may vary. If the physical size $L$ of the bubble is small, $L \leq cH^{-1}(z_i) \Delta \nu_{\text{obs}}/\nu_{\text{obs}}$, then $\Delta r \approx L$. On the other hand, if $L \geq cH^{-1}(z_i) \Delta \nu_{\text{obs}}/\nu_{\text{obs}}$ and the bubble has split from the Hubble flow, then $\Delta r = L$. But if the bubble is comoving with the Hubble flow, then $\Delta r = cH^{-1}(z_i) \Delta \nu_{\text{obs}}/\nu_{\text{obs}}$ (this length, at $z = 7$, corresponds to $\sim 12$ kpc). Note that, unless otherwise stated, distances will always be in physical units (and not comoving). In this work we assume that bubbles are not gravitationally bound, and for simplicity, we further assume that they expand with the Hubble flow, so that $\Delta r = \min[L, cH^{-1}(z_i) \Delta \nu_{\text{obs}}/\nu_{\text{obs}}]$. Hence, in this case,

$$y \approx \frac{j_\nu}{B_\nu(T_0)} \min \left[ L, cH^{-1}(z_i) \frac{\Delta \nu_{\text{obs}}}{\nu_{\text{obs}}} \right].$$  

(4)

Depending on the angular and spectral resolution of a specific instrument, the $y$-distortion along a given line of sight, and at a given frequency, may have to be convolved with the respective response functions (clearly, such smoothing can reduce the fluctuations). In what follows, we adopt a model instrumental beam (or point-spread function [PSF]) response function $V_{\text{PSF}}(\hat{n})$ and a frequency response function $\Phi(\nu)$. Both will be taken to have accurately determined Gaussian shapes, so that $\int d\hat{n} V_{\text{PSF}}(\hat{n}) = \int d\nu \Phi(\nu) = 1$, where $\hat{n}$ denotes the unit vector.

We assume generically that each $O_1$ bubble is associated with a halo of some mass $M$, with an abundance given by the Sheth-Tormen (Sheth & Tormen 2004) mass function $dn/dM$, and that the $O_1$ atoms are distributed in a bubble around the halo with a spherically symmetric density profile $W_b(r)$; in this case, a second convolution, over the bubble density profile, is necessary.

If we neglect the peculiar velocities and internal motions of the bubbles, the effective distortion along the line of sight in direction $\hat{n}$ and at frequency $\nu_{\text{obs}}$ is

$$y_{\text{eff}}(\nu_{\text{obs}}, \hat{n}) = \int_0^\infty \int_{psf} d\nu d\hat{n} \Phi(\nu) V_{\text{PSF}}(\hat{n}) \int_{r_{\nu} + \Delta r/2}^{r_{\nu} - \Delta r/2} dr \times \int dy dM \frac{dn}{dM}(y, M) W_b(y - r) \frac{j_\nu}{B_\nu(T_0)},$$  

(5)

where $r = r_{\nu}$. The distance $r$ is centered at the resonant value $r_{\nu} \equiv r(z_i)$ and $\Delta r$ corresponds to the effective length defined above. The integral over $M$ counts $O_1$ bubbles around halos of various masses (above some minimum halo mass $M_{\text{min}}$, as discussed further in § 5 below), whereas the integrals over $\nu$ and $\hat{n}$ describe the cosmological volume probed by the PSF and the spectral response of the detector. Let us change variables from frequencies and angles to spatial coordinates, and define the three-dimensional instrumental PSF:

$$B(x - r_{\nu} \equiv \Phi(\nu)(\nu_{\text{obs}}/r_{\nu}^2)(d\nu/dz)(\nu)/V_{\text{PSF}}(\hat{n}).$$  

Further defining $y(M, z) \equiv \{j_{\nu}/B_{\nu}(T_0)\} \Delta r$, the distortion generated by a bubble around a single halo of mass $M$ at redshift $z$, we can rewrite equation (5) more transparently as

$$y_{\text{eff}}(\nu_{\text{obs}}, \hat{n}) \approx \int dM \left[ B \ast \left( \frac{dn}{dM} * W_b \right) \right] \hat{y},$$  

(6)

where an asterisk denotes convolution. In what follows, the twice-convolved halo number density is denoted by $\tilde{n}$, i.e., $\tilde{n}(M, r_{\nu}) \equiv \{B \ast (dn/dM \ast W_b)(M, r_{\nu})\}$.

Finally, the physical collision between two metal-rich bubbles will lead to a complicated thermodynamical interaction. Rather than modeling this process, the above prescription assumes a linear addition of the $O_1$ density (or $y$-distortion) from two physically overlapping metal bubbles. This simple assumption at least captures the enhancement of the metallicity in the overlap regions. We note that this is different from the case of merging H II bubbles during reionization, when mergers conserve volume and result in the expansion of the joint bubble.

4. The Effect of Clustering

We next compute the second-order moments (correlation function and power spectrum) of the distortion field generated by $O_1$ at 63.2 $\mu$m during reionization. Throughout this discussion, we consider a fixed redshift and suppress the frequency dependence of $y_{\text{eff}}(\hat{n})$. The angular correlation function can be written as

$$\langle y_{\text{eff}}(\hat{n}_1) y_{\text{eff}}(\hat{n}_2) \rangle = \int \int dM_1 dM_2 \tilde{n}(\hat{n}_1) \tilde{n}(\hat{n}_2) \hat{y}_1 \hat{y}_2 = \int \int dM_1 dM_2 \tilde{n} \tilde{n} \hat{y}_1 \hat{y}_2 [1 + \xi_{\text{hh}}(r_1, M_1, r_2, M_2)].$$  

(7)

In this equation, $\tilde{n}$ denotes the average $\tilde{n}$. The correlation function $\xi_{\text{hh}}(r_1, M_1, r_2, M_2)$ corresponds to the halo-halo correlation function convolved with both the window function of the profile of the $O_1$ distribution ($W_b$) and the window function of the experiment ($B$). Keeping in mind that the Fourier counterpart of $\xi_{\text{hh}}$ is the matter power spectrum times the square of the bias factor $b(M, z)$ computed in Sheth & Tormen (1999), equation (7) becomes

$$\langle y_{\text{eff}}(\hat{n}_1) y_{\text{eff}}(\hat{n}_2) \rangle = \int \int dM_1 dM_2 \tilde{n} \tilde{n} \hat{y}_1 \hat{y}_2 [1 + \xi_{\text{hh}}(r_1, M_1, r_2, M_2)] \times \int \frac{dk}{(2\pi)^3} P_{\text{m}}(k, z_e) |\mathcal{E}_k|^2 |W_{b, h}|^2 \exp[-i \mathbf{k} \cdot (r_1 - r_2)],$$  

(8)

where $P_{\text{m}}(k, z_e)$ the matter power spectrum at redshift $z_e$ (note that we have dropped a constant $[k = 0]$ term, and hence we are looking at the departure of the correlation function from its mean value) and $W_{b, h}$ and $B_k$ the Fourier counterparts of $W_b$ and $B$, respectively. Since we are assuming that every halo is producing a bubble, and the dominant signal will be due to the clustering of different bubbles, we can neglect nonlinear corrections to the matter power spectrum. Such nonlinear corrections would boost the clustering signals we predict below, but $^5$ The integral $\int dx B(x - r_{\nu}) = 1$ is normalized to unity in the volume centered at $r_{\nu}$. 
only on comoving scales of $\leq 0.1$ Mpc at $z \approx 6$ (e.g., Iliev et al. 2003), which is well below the scale at which the O\textsc{i} clustering signal peaks. The integral over $k$ can be split into a transverse ($k_\perp$) and a parallel ($k_z$) component along the line of sight. We assume that both $B_\perp$ and $W_{b,\perp}$ can be factorized as $B_\perp = (B_k, B_{h,k})$ and $W_{b,\perp} = (W_{h,k} W_{b,k})$. For a Gaussian instrumental response, we can write $B_k = \exp[-k_z^2 (r_\perp \sigma_r)^2/2]$ and $B_k = \exp[-k_z^2 (dr/dz/\sigma_r/\sigma_{\text{obs}})^2/2]$, with $\sigma_r$ related to the width of the angular PSF and $\sigma_{\text{obs}}$ to the width of $\Phi(v)$. For simplicity, the bubble is taken to have a Gaussian profile and a volume equal to $(2\pi)^3/2 L^3$. After this decomposition, the correlation function reads

$$
\langle n_{\text{eff}}(n_1) n_{\text{eff}}(n_2) \rangle = \int \frac{dk}{2\pi} \exp[-ik_\perp (r_1 - r_2)] \int dM_1 dM_2 n_1 n_2 b_1 b_2 \times |B_{k,\perp} W_{b,\perp}|^2 \int \frac{dk_z}{2\pi} |B_{k,\perp} W_{b,\perp}|^2 P_m(k_\perp, k_z, z_\perp). \tag{9}
$$

Finally, by noting that in the flat-sky approximation $k_\perp \approx l/r_s$, we obtain the angular power spectrum:

$$
C_l = \frac{1}{s} \int \int dM_1 dM_2 n_1 n_2 b_1 b_2 |B_{k,\perp} W_{b,\perp}|^2 \int \frac{dk_z}{2\pi} |B_{k,\perp} W_{b,\perp}|^2 P_m(l/r_s, k_z, z_\perp). \tag{10}
$$

Let us briefly examine the behavior of the $C_l$ values. We first recall that the Fourier window function of a bubble is proportional to the bubble volume ($\propto L^3$), and that the oxygen number density in bubbles is proportional to $n_0/(N L^3)$, with $N$ the average bubble number density, $n_0$, the global mean oxygen number density at a given redshift, and $L$ the typical bubble size. We also note that the bias factor $b$ almost cancels the redshift dependence of the growth factor of perturbations (e.g., Oh et al. 2003), so the scaling of the maximum of $f^2 C_l$ is $(f C_l)_{\text{max}} \propto G^2 n_0^2 (\Delta r)^2$, with $G$ a frequency-dependent function accounting for the efficiency of the pumping process and $\Delta r$ the effective length in equation (3). The term $G$ dominates the overall redshift dependence of the power-spectrum amplitude. If $L$ is smaller than $cH^{-1}(z)\Delta r_{0,10}/v_{10}$, then $(f C_l)_{\text{max}} \propto G^2 n_0 L^2$, but otherwise, $(f C_l)_{\text{max}} \propto G^2 n_0^2$, i.e., the band power spectrum is independent of the bubble size (except that the bubble Fourier window function suppresses sources at scales smaller than the bubble size). In either case, note that $(f^2 C_l)_{\text{max}}$ is proportional to the square of the oxygen metallicity. However, if the O\textsc{i} abundance within the bubbles is so high that they become optically thick at $\lambda_{20}$, then $y$ will reach its plateau, $y \sim \tilde{G}$, and $(f^2 C_l)_{\text{max}} \propto G^2 (N L^3)^2$, i.e., $(f^2 C_l)_{\text{max}}$ will be proportional to the square of the number density of bubbles times their cross-sectional area, assuming that the frequency resolution of the instrument corresponds to a radial distance that does not exceed the typical bubble size. Otherwise, the scaling would be $(f^2 C_l)_{\text{max}} \propto G^2 (N L^3)^2$. The discrete nature of the source distribution means that in addition to the correlation term computed above, there will be a Poisson term. The Poisson contribution can be computed by considering the limiting case of equation (10) for a random source distribution and an experiment of infinite angular and spectral resolution,

$$
C_l \approx y^2 (\tilde{N} L^3) \Delta \Omega_b, \tag{11}
$$

with $y$ the average bubble distortion, $\tilde{N} L^3$ the volume fraction occupied by bubbles, and $\Delta \Omega_b$ their typical angular size. In practice, we find that the shot-noise contribution is nearly always subdominant in the results we present below.

5. Toy Model for Metal Distribution

To explore the detectability of the angular power spectrum of the O\textsc{i}–induced CMB distortion, we need a model for the spatial distribution of metals in the IGM. Following the discussion in §§3 and 4 above, this means we need to specify the profile $W_q(r)$ around a halo of mass $M$ at each redshift $z$.

Our basic simplifying assumption is that metals are confined in bubbles that had previously expanded into the IGM, but that these bubbles have settled to follow the Hubble flow. We assume that when a halo is formed, a metal-polluting wind is launched from the galaxy at the center of the halo, producing a bubble that subsequently expands at a constant velocity for the age of the halo. In principle, halos of the same mass that exist at a fixed redshift $z$ can have a distribution of ages. For simplicity, however, we assign a fixed age to halos of mass $M$ at redshift $z$, corresponding to the average redshift $z_f$ at which such halos have assembled 65% of their mass (see, e.g., eq. [2.26] in Lacey & Cole 1993). At redshift $z$, the radius of the bubble around a halo of mass $M$ is therefore given by

$$
L(M, z) = 0.1 r_{\text{vir}}(M, z_f) + (v_z^2 - v_{\text{esc}}^2)^{1/2} [t(z) - t(z_f)], \tag{12}
$$

with $r_{\text{vir}}$ the virial radius of the parent halo and $t(z)$ the cosmic time at redshift $z$. The first term on the right-hand side represents a rough estimate for the size of the galaxy in the halo (i.e., the region producing the metals). The velocity $v_z$, taken to be a constant in the range of 50–1500 km s$^{-1}$, represents the wind velocity “in vacuum.” We then subtract $v_{\text{esc}}$, corresponding to the

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig2.png}
\caption{Radius of metal bubbles around halos of mass $M$ produced in our fiducial toy model, assuming metals expand over the typical age of the halo at a constant speed of $(v_z^2 - v_{\text{esc}}^2)^{1/2}$. Here $v_z$ is a wind velocity equal to 150 km s$^{-1}$, and $v_{\text{esc}}$ is the escape velocity from the halo. The solid curve shows the radii at $z = 6.5$, and the dashed curve at $z = 7.5$. The sharp feature at $M \approx 10^9 M_\odot$ corresponds to the limit where metals are trapped inside the halo potential well ($r_f \approx r_{\text{esc}}$). The bubble radius for more massive halos is assumed to be a fixed fraction (here taken to be 10%) of the virial radius, and so it scales as $M^{1/3}$.}
\end{figure}
escape velocity from the parent halo at a distance of \(v_{\text{esc}}(M, z_f)\), from this wind velocity, to take into account the fact that metals lose energy as they travel out of the gravitational potential well. Since we assume that metals have no peculiar velocities, the bubbles start comoving with the Hubble flow by the time they are observed. Equation (12) assumes further that star formation takes place instantly when the halo forms. In Figure 2 we show the bubble sizes predicted under these assumptions.

For simplicity, we normalize the total amount of oxygen mass in all bubbles at a given redshift by specifying the global average metallicity of the IGM at each redshift, i.e.,

\[
Z_{av}(z) = \int_{M_{\text{min}}} M dM \frac{d n}{d M}(M, z) Z(M, z) V_b(M, z),
\]

where \(Z(M, z)\) is the metallicity within a bubble at redshift \(z\) originated in a halo of mass \(M\) and with volume \(V_b = 4\pi/3L^3\), and \(d n/d M\) is the halo mass function. For halos with \(v_b > v_{\text{esc}}\), we further assume that, at a given redshift, the oxygen abundance within the bubble is proportional to the mass of the parent halo. Provided that the bubble radii are practically constant in this mass range, this is roughly equivalent to assuming that the star formation rate scales linearly with halo mass and is constant throughout the age of the halo. For halos with \(v_b < v_{\text{esc}}\), we assume that the oxygen density is constant (since the metals produced in these halos are assumed to be trapped in a volume \(V_b \propto M\)); however, these larger halos contribute negligibly to the clustering signal below, and our results are insensitive to this assumption.

In Figure 3 we show the evolution of the global average (dashed lines) and bubble-volume-weighted (solid lines) metallicities versus redshift. The dotted lines show the fraction of the total volume contained inside the bubbles (note that metal-rich bubbles can overlap, and this fraction can exceed unity). In all panels, the global average oxygen abundance is scaled to \(10^{-2.5} Z_0\) at each redshift. This value is not unrealistic, as there are observational constraints on the cosmic metallicity indicating that \(Z \sim 10^{-3} Z_0\) at \(z > 4\) (Schaye et al. 2003), and oxygen from first-generation stars is likely to be somewhat overabundant relative to the solar value (Meynet et al. 2006). The highest redshift, \(z = 12\), shown in the figures represents the earliest epoch when this floor metallicity may have been established (e.g., Haiman & Loeb 1997). Note that since the global metallicity is rescaled at each redshift by hand, the curves in Figure 3 should be interpreted as independent models of the metal distribution at each redshift, and not as evolutionary models. In order to illustrate the metal distribution under different assumptions, in Figure 3a we assume \(M_{\text{min}} = 5 \times 10^5 M_\odot\), and \(v_b = 50 \text{ km s}^{-1}\); Figure 3b is the same as Figure 3a, except with \(v_b = 150 \text{ km s}^{-1}\), while in Figure 3c we assume \(v_b = 150 \text{ km s}^{-1}\) and \(M_{\text{min}} = 5 \times 10^7 M_\odot\).

The enrichment scenarios considered in this work are comparable with those presented in previous studies. Indeed, in Furlanetto & Loeb (2003) the typical bubbles sizes at \(z \sim 6\) are a few tens of kiloparsecs, in agreement with our estimates, and in addition, after looking at the simulations of Oppenheimer et al. (2007) we have found similar enrichment levels in the relevant redshift ranges (\(z \in [5, 10]\)). Only in Scannapieco et al. (2002) do the bubble filling factors tend to be slightly smaller than in our case, since, motivated by WMAP observations, we are considering early reionization models.

6. RESULTS AND DISCUSSION

The angular power spectrum (correlation term plus shot-noise term) of the CMB distortion produced by O I pumping is shown in Figure 4. We have adopted an angular resolution of \(\theta^p\) and a relative spectral resolution of \(\Delta v_{\nu}/v = 10^{-4}\). The angular resolution corresponds to \(\ell \sim 7 \times 10^2\), and therefore should resolve all bubbles under consideration. The frequency resolution corresponds at \(z = 6.5\) to roughly 60 kpc in physical units. In all three panels we consider variations from our adopted fiducial model given by \(z = 6.5\) (observable at 630 GHz), \(I_{\nu}^{p} = 6 \times 10^{-20} \text{ erg cm}^{-2} \text{ s}^{-1} \text{ Hz}^{-1} \text{ sr}^{-1}\), \(M_{\text{min}} = 5 \times 10^5 M_\odot\), and \(v_b = 150 \text{ km s}^{-1}\), corresponding to \(L \sim 60 \text{ kpc}\) at that redshift. Figure 4a shows the dependence of the signal on bubble expansion/volume/bubble size: 150 (solid curve), 60 (dashed curve), and 200 (dotted curve) \(\text{km s}^{-1}\), which correspond to \(L \sim 60, 23\), and 80 kpc, respectively. If \(v_b = 60 \text{ km s}^{-1}\) (\(L \sim 23 \text{ kpc}\)), then \(L < cH^{-1}(z)\Delta v_{\nu}/v_0\), and the bubbles remain optically thin (\(y\) does not reach its maximum value), so that \((L^2 C_n)_{\text{max}} \propto G^2 L^2\). On the other hand, for \(v_b \sim 150 \text{ km s}^{-1}\) (\(L \sim 60 \text{ kpc}\)), we are in the regime of \(L \sim cH^{-1}(z)\Delta v_{\nu}/v_0\) and \((L^2 C_n)_{\text{max}} \propto G^2\). Therefore,
if $v_\text{b} = 200$ km s$^{-1}$, the amplitude of the spectrum should not change noticeably. However, the bubble size ($L \approx 80$ kpc) is slightly larger than in our fiducial model ($L \approx 60$ kpc), and when comparing both cases one can clearly see the effect of the bubble window function suppressing power at small scales. Therefore, $(\ell^2 C_\ell)_{\text{max}} \propto (v_\text{b}/150$ km s$^{-1})^2$ if $v_\text{b} < 150$ km s$^{-1}$. In Figure 4b we find a strong dependence of the power on the redshift: for the range of $z_r = 6.5, 7,$ and $7.5$ (top to bottom), the amplitude of the power spectrum varies by nearly 2 orders of magnitude. This steep dependence arises because the efficiency of the pumping is power spectrum varies by nearly 2 orders of magnitude. This dependence on the minimum halo mass: $M_\text{min} = 5 \times 10^5$ (solid line), $5 \times 10^6$ (dashed line), and $7.5 \times 10^6$ M$_\odot$ (dotted line). The distortion $y$ becomes saturated for $M_\text{min} \gtrsim 5 \times 10^6 M_\odot$. For $5 \times 10^5 M_\odot < M_\text{min} < 5 \times 10^6 M_\odot$ we have the approximate scaling $(\ell^2 C_\ell)_{\text{max}} \propto y^2$, which is slightly distorted by the presence of the bias factor enhancing the dashed curve over the solid one. For $M_\text{min} = 7.5 \times 10^6 M_\odot$ we are well into the optically thick regime [$C_\ell \propto (NL^3 y)^2$], and the typical bubble sizes have shrunk dramatically (see Fig. 2; this translates into a clear drop in the power (thinnest curve)).

In this paper we have computed the angular fluctuation of the effect we presented in an earlier work (Paper I): the CMB spectral distortion induced by the O i Balmer-$\alpha$ line. Since the oxygen distribution in space is associated with the most overdense regions hosting the first star-forming activity at the end of the Dark Ages, there must be fluctuations in the angular pattern of the O i-induced y-distortion. We have considered the clustering properties of the signal and computed its angular power spectrum using a toy model for the spectrum peaks. In principle, the typical bubble clustering length could also be detected in the radial direction if the observing instrument had enough spectral resolution: in such a case, observations at slightly different frequencies would give rise to correlated maps, since they would be probing overlapping shells centered at similar redshifts. The combination of maps obtained at different frequencies could also improve the signal-to-noise ratio of the final detection.

Because at the frequencies of interest the signal will be dominated by infrared emission from dusty galaxies, it is important to be able to remove them from the observed map. One way this can be achieved is by obtaining high angular resolution images of the sky at the frequencies of interest, so that Olber’s paradox is avoided and individual galaxies can be clipped from the maps. ALMA may be an excellent instrument for this purpose because of its high sensitivity and angular resolution (of the order of 0.5$''$). The drawback is that the field of view of ALMA is small ($\leq 1'$). However, since our signal peaks at scales of approximately tens of arcseconds and decays slowly at smaller scales, this small field of view should not be a fundamental limitation. The signal could also be within reach of the sensitivity of forthcoming detectors (e.g., SCUBA2; Holland et al. 2006) and of planned single-dish experiments such as CCAT. However, current small-scale CMB experiments, with a typical sensitivity of $\sim 1 \mu$K at arcminute scales, are still rather far from imposing interesting constraints on this effect.

7. CONCLUSIONS

In this paper we have computed the angular fluctuation of the effect we presented in an earlier work (Paper I): the CMB spectral distortion induced by the O i Balmer-$\alpha$ line. Since the oxygen distribution in space is associated with the most overdense regions hosting the first star-forming activity at the end of the Dark Ages, there must be fluctuations in the angular pattern of the O i-induced y-distortion. We have considered the clustering properties of the signal and computed its angular power spectrum using a toy model for the
metal distribution. We find that the predicted signal is small, but that for certain enrichment models it could be detected using ultradeep observations with sensitive instruments such as ALMA, SCUBA2, and CCAT. Even in the event of a nondetection, future infrared observations could place interesting limits on the metal enrichment of the universe during the Dark Ages. This could open the possibility of measuring the metal enrichment of the universe before it was reionized. OI observations would also be complementary to HI 21 cm measurements that will map out in the near future the distribution of neutral H during the epoch of reionization.

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