The Astrophysical Journal, 664:1–7, 2007 July 20

HYDROGEN 2*p*–2*s* TRANSITION: SIGNALS FROM THE EPOCHS OF RECOMBINATION AND REIONIZATION

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Received 2007 February 27; accepted 2007 April 2

ABSTRACT

We propose a method to study the epoch of reionization based on the possible observation of 2*p*–2*s* fine-structure lines from the neutral hydrogen outside the cosmological H II regions enveloping QSOs and other ionizing sources in the reionization era. We show that for parameters typical of luminous sources observed at $z \simeq 6.3$, the strength of this signal, which is proportional to the H I fraction, has a brightness temperature $\simeq 20$ μK for a fully neutral medium. The fine-structure line from this redshift is observable at $\nu \simeq 1$ GHz, and we discuss prospects for the detection with several operational and future radio telescopes. We also compute the characteristics of this signal from the epoch of recombination; the peak brightness is expected to be $\simeq 100$ μK, and this signal appears in the frequency range 5–10 MHz. The signal from the recombination era is nearly impossible to detect owing to the extreme brightness of the Galactic emission at these frequencies.

Subject headings: cosmic microwave background — line: formation — radiative transfer — radio lines: general

1. INTRODUCTION

Even though the existence of hydrogen fine-structure lines and their explanation using Dirac's atomic theory has been known for close to a century, a hydrogen fine-structure line has never been detected from an astrophysical object. An interesting hydrogen fine-structure line is the 2*p*–2*s* transition. The main difficulty in detecting this line is that the line strength is proportional to the population of either the 2*p* or 2*s* states, which, being excited states, are not so readily populated in most astrophysical circumstances. Moreover, the line width of the excited 2*p* state, which is determined by its decay time, is large (99.8 MHz), making the detection of the fine-structure line a difficult observation. One astrophysical setting where the feasibility of detecting such a line has been studied is the interstellar H II regions (see, e.g., Dennison et al. 2005 and references therein; Ershov 1987), in which the excited levels are populated by recombination.

Here we consider two cosmological settings in which the excited levels are populated by either recombination or pumping by Lyα photons from an external source.

The recombination epoch.—The universe makes a transition from a fully ionized to an almost fully neutral medium at $z \simeq 1089$ (Spergel et al. 2007; for details see, e.g., Peebles 1993 and references therein). During this era, as the density and temperature of the universe drops, recombination is stalled, owing to a high Lyα radiation density, and progresses by either the depopulation of the 2*p* state, owing to redshifting of the photons out of the line width, or the two-photon decay of the 2*s* state. This results in a significant 2*s* and 2*p* level population during the recombination era.

The reionization epoch.—Recent observations suggest that the universe made a transition from nearly fully neutral to fully ionized within the redshift range $6 \lesssim z \lesssim 15$ (Page et al. 2007; White et al. 2003; Fan et al. 2002; Djorgovski et al. 2001; Becker et al. 2001). It is widely believed that this "reionization" was achieved by the percolation of individual H II regions around the sources of reionization. The nature of these sources is not well understood; they might be Population III stars, active galactic nuclei, or star-forming galaxies. During this epoch, a signal from the 2*p*–2*s* fine-structure transition might originate either from within the cosmological H II regions or from the almost fully neutral medium surrounding the H II region. The level population of the first excited state in the former case would be largely determined by recombinations and in the latter case by Lyα photons from the central source. We show below that for most cases of interest, the fine-structure line from within the cosmological H II region might be negligible as compared to the signal from the regions immediately surrounding the H II region.

Throughout this work we adopt the currently favored ΛCDM model, spatially flat with $\Omega_m = 0.3$ and $\Omega_\Lambda = 0.7$ (Spergel et al. 2007; Riess et al. 2004; Perlmutter et al. 1999), with $\Omega_b h^2 = 0.022$ (Spergel et al. 2007; Tytler et al. 2000) and $h = 0.7$ (Freedman et al. 2001).

2. FINE-STRUCTURE LINES FROM THE REIONIZATION EPOCH

Subsequent to the recombination of the primeval baryon gas at redshift $z \simeq 1089$ (Spergel et al. 2007) and the transformation of the gas to an almost completely neutral state, it is believed that the gas was reionized during epochs corresponding to the redshift range $6 \lesssim z \lesssim 15$. *Wilkinson Microwave Anisotropy Probe* (*WMAP*) measurements of cosmic microwave background radiation (CMBR) anisotropy in total intensity and polarization have been used to infer that the baryons were likely neutral at redshifts $z \gtrsim 12$–15; however, the detection of CMB polarization anisotropy requires substantial ionization by about $z \simeq 11$ (Page et al. 2007). Observationally, the Gunn-Peterson (GP) test shows that the universe is highly ionized at redshifts lower than $z \simeq 5.5$; the detection of GP absorption at greater redshifts suggests that the neutral fraction of the intergalactic hydrogen gas rises to at least 10^{-3} in the redshift range $5.5 \lesssim z \lesssim 6$ and that reionization was not complete until about $z \simeq 6$ (White et al. 2003; Fan et al. 2002; Djorgovski et al. 2001; Becker et al. 2001). However, from the GP test alone it is not possible to infer the neutral fraction of the medium; it only gives a rather weak lower limit of $\simeq 10^{-3}$ on the neutral fraction of the universe for $z \gtrsim 6$. From other considerations it is possible to put more stringent bounds on the neutral fraction; for example, Wyithe & Loeb (2004) obtain a lower limit of 0.1 on the neutral fraction of the universe at $z \simeq 6.3$ (see also Mesinger & Haiman 2004).

Our understanding of the nature of the sources that caused the reionization is far from complete. The transition from an almost completely neutral gas to a highly ionized gas during redshifts $6 \lesssim z \lesssim 15$ is a key problem in modern cosmology and considerable theoretical and experimental efforts are currently directed at this unsolved problem. Here we propose a new method, based on the $2p$–$2s$ fine-structure transition, for determining the evolution of the neutral fraction of the intergalactic medium within this epoch.

Owing to fine-structure splitting, the two possible transitions between the $2s$ and the $2p$ states are $2p_{1/2}$–$2s_{1/2}$ at a frequency $\simeq 1058$ MHz that has an Einstein A-coefficient 1.6×10^{-9} s^{-1} and $2p_{3/2}$–$2s_{1/2}$ at a frequency $\nu_{ps} \simeq 9911$ MHz that has an Einstein A-coefficient 8.78×10^{-7} s^{-1}. The Einstein A-coefficient for the latter transition is more than an order of magnitude greater than the former; therefore, in this work we consider only the $2p_{3/2}$–$2s_{1/2}$ transition and hereafter refer to this specific transition as simply the $2p$–$2s$ transition.

The ionizing UV photons from sources in the reionization era create "Stromgren spheres." Whereas the gas in the cosmological H II regions are highly ionized by the photons, the ionization level of the gas beyond the Stromgren spheres is determined by the history of the gas, the density, and the mean specific intensity of the background ionizing photons, which includes both the photons diffusing out of the Stromgren spheres as well as the background radiation field. In this work we assume that the ionizing sources at these high redshifts are AGNs, and in illustrative examples adopt parameters of a few QSOs that have been observed at $z \simeq 6$. The photons at the Lyα transition frequency (here and throughout, unless otherwise specified, we continue to refer to frequencies between Lyα and Lyman-limit as "Lyα") from a high-redshift QSO escape the mostly ionized Stromgren sphere and are strongly scattered and absorbed in the medium beyond. The population of the $2p$ level in this region is determined by (1) the intensity of Lyα photons from the central source, (2) recombination rate of free electrons, (3) absorption of CMBR photons by electrons in the $2s$ state (it is assumed here and throughout this work that the only radio source at high redshifts is the CMBR), and (4) collisional transition from the $2s$ state. The $2s$ state is populated via (1) the recombination rate of free electrons, (2) collisional transfer of atoms from the $2p$ state, (3) the spontaneous decay of the $2p$ state, and (4) transition from the $2p$ state stimulated by CMBR photons. In addition, the absorption of photons from the central sources, with energy equal to or in excess of the Lyβ transition, would result in electronic transitions to the second excited state, which could be followed by spontaneous decay to the $2s$ state. (It might be pointed out here that both the $2s$ and $2p$ states could also be populated by atoms cascading from excited states with $n > 3$. In particular, all photons absorbed from $1s$ states to any excited state can directly deexcite to the $2s$ level. However, the rate of transition from $1s$ to any excited state is roughly $\propto 1/n^3$ [e.g., Rybicki & Lightman 1979], and therefore, we include only the most dominant transition in each case.)

The population of the ground state is denoted by n_{1s}. We denote the level populations of the two states ($2s_{1/2}$ and $2p_{3/2}$) by the number densities n_{2s} and n_{2p}; these may be solved for, respectively, from the two equations of detailed balance

$$f\alpha_B n_i^2 + cB_{2p2s}n_{2p}n_{\rm CMBR} + C_{ps}n_i n_{2p} + A_{2p2s}n_{2p} + cn_{1s}p_{32}\int B_{13,\beta}\phi_{13}(\nu)n_\alpha(\nu)d\nu = A_{2s1s}n_{2s} + C_{sp}n_i n_{2s} + cB_{2s2p}n_{2s}n_{\rm CMBR}, \quad (1)$$

$$(1-2f)\alpha_B n_i^2 + cB_{2s2p}n_{2s}n_{\rm CMBR} + C_{sp}n_i n_{2s} + cn_{1s}\int B_{12,\alpha}\phi_{12}(\nu)n_\alpha(\nu)d\nu = A_{2p1s}n_{2p} + C_{ps}n_i n_{2p} + cB_{2p2s}n_{2p}n_{\rm CMBR}, \quad (2)$$

where f is the fraction of all the atoms that recombine to the $2s$ state. In equilibrium $f = 1/3$ as the $n = 2$ state splits into three doublets: $2p_{1/2}$, $2p_{3/2}$, and $2s_{1/2}$; α_B is the recombination coefficient and n_i is the density of the ionized gas; $B_{2p2s} = B_{2s2p} = c^2/(8\pi\nu_{ps}^3)A_{2p2s}$ is the Einstein B-coefficient for the $2p_{3/2}$–$2s_{1/2}$ transition in terms of the corresponding Einstein A-coefficient A_{2p2s} (note that the two B-coefficients are equal as the two states have the same degeneracy); $C_{ps} = C_{sp} = 5.31 \times 10^{-4}$ cm^3 s^{-1} is the rate coefficient of transition owing to collisions with electrons; $n_{\rm CMBR}$ is the number density of CMBR photons within the transition line width; $n_\alpha(\nu)$ is the number density (per unit frequency) of photons with frequency equal to or larger than the Lyα frequency (and smaller than the Lyman-limit frequency) at any location; ϕ_{13} and ϕ_{12} are, respectively, the line profiles corresponding to the Lyβ and Lyα transitions, and p_{32} is the probability for the electron transition to the $2s$ state following excitation to $n = 3$ via absorption of a Lyβ photon. We have not included the induced Lyα transition, because the number density of atoms in the $2p$ state is negligible as compared to that in the $1s$ state; p_{32} is the probability that an atom in the third excited state ($3p$) will decay to the $2s$ state; A_{2s1s} is the Einstein A-coefficient corresponding to the two-photon decay of the $2s$ state. Other symbols have their usual meanings.

Owing to the fact that the mostly neutral medium in the vicinity of the cosmological Stromgren spheres is optically thick to Lyα scattering, the Lyα photons from the decay of the $2p$ state are strongly scattered by the gas. Therefore, $n_\alpha(\nu)$ will contain contributions from both the Lyα photons from the central source as well as the Lyα photons that arise from recombinations outside the Stromgren sphere and are multiply scattered therein, $n_\alpha(\nu) = n_\alpha^{\rm src}(\nu) + n_\alpha^{\rm rec}(\nu)$. We neglect the multiply scattered Lyα photons from the Stromgren sphere that have been reprocessed via recombination within the Stromgren sphere, because these would be redshifted redward of the Lyα line before encountering the boundary of the Stromgren sphere. The scattering of recombination photons in an optically thick, expanding medium is a complex problem (Field 1959; Rybicki & Dell'Antonio 1994). One of its first applications was to study the recombination of primeval plasma (Peebles 1968; Zeldovich et al. 1969). In these analyses it was implicitly assumed that apart from two-photon decay, in an expanding universe the dominant effect that results in resonant photons ceasing interaction with the gas, and leaving the system, is its redshifting out of the line profile. The effect of scattering off the moving atoms was deemed to be either negligible or at best comparable. This assumption has been borne out by more recent detailed analysis that has taken into account the effect of scattering on the photon escape (Krolik 1990). Taking only the redshift as the main agent of photon escape, it can be shown that the net effect of the scattering of a resonance photon before it drops out of consideration is to reduce the decay time of the $2p$ state from $A_{21} = 6.2 \times 10^8$ s^{-1} to $A_{21}/\tau_{\rm GP}$ (Zeldovich et al. 1969; for more recent work see, e.g., Chluba et al. [2007], Seager et al. [1999], and references therein; we give a concise derivation in the Appendix). In this expression, the GP optical depth $\tau_{\rm GP} = [3/(8\pi H)]A_{2p1s}\lambda_\alpha^3 n_{1s}$.

Similar complications exist in computing p_{32}, the probability that an atom in the third excited state will decay spontaneously to the $2s$ state, in an optically thick medium. In optically thin media, $p_{32} = A_{32}/(A_{32} + A_{31})$; on the other hand, in an optically thick medium, we would need to take into account the "trapping" of the Lyβ photon owing to resonant scattering. The effect of this scattering in an optically thick medium would be that the fraction of photons that decay directly to the ground state are reabsorbed "locally" to the third excited state, and therefore, all photons absorbed to the third excited state result in an Hα photon and an atom in the $2s$ state. This means that the appropriate value of p_{32} is close to unity in an optically thick medium; in this work we assume $p_{32} = 1$.

The astrophysical setting in which we seek solutions to the algebraic equations above is cosmological H II regions at high redshift. In particular, we are interested in the signal from the neutral region surrounding the cosmological H II region. For a fully neutral intergalactic medium (IGM) at $z \simeq 6.5$, $\tau_{\rm GP} \simeq 6 \times 10^5$. If we adopt spectral luminosities corresponding to QSOs observed at these high redshifts, it may be shown that the populating of the $2p$ state via direct recombinations from the free-free state, collisional transfer from the $2s$ state, and upward transitions from the $2s$ to $2p$ state arising from absorption of background CMBR photons may all be neglected. (The relevant parameters are $n_i \simeq n_b \simeq 2.8[1+z]^3$ cm^{-3} in the H II region surrounding the sources, with n_i expected to be much smaller in the neighboring, mostly neutral medium; the number density of CMBR photons that might cause a $2p$–$2s$ transition is $\simeq 5[1+z]$ cm^{-3}; the number density of Lyα photons from the central source, assuming luminosities typical of SDSS quasars at $z \simeq 6.5$ [more details in § 4], is $n_\alpha \simeq 10^{-4}$ cm^{-3}. First, for these parameters the dominant process that populates the excited state is the pumping by Lyα photons. Second, it may be readily verified that for these plausible values for the parameters, the signal expected from the H II region surrounding the central source is negligible as compared to the signal from the surrounding neutral region.) Given that the Lyα flux from the central QSO is the dominant causative factor for populating the $2p$ state, the two equations (eqs. [1] and [2]) that determine the level populations are simplified. The number density of atoms in the $2p$ state is given approximately by

$$n_{2p} \simeq n_{1s}\Gamma_\alpha \tau_{\rm GP}/A_{2p1s}, \tag{3}$$

where $n_{1s} = f_{\rm neu} n_b$, with $f_{\rm neu}$ denoting the neutral fraction and $n_b \simeq 2.7 \times 10^{-7}(1+z)^3$ cm^{-3} being the number density of baryons in the IGM; $\Gamma_\alpha = \int B_{2p1s}\phi_{12}(\nu)n_\alpha^{\rm src}(\nu)$ (in s^{-1}) is the transition rate to the $2p$ state owing to the Lyα photons from the central source, where $n_\alpha^{\rm src}(\nu)$ is the number density of Lyα photons from the central source alone.

Similarly, the dominant process that determines the population of the $2s$ state is the absorption of Lyβ photons (for details, see the discussion above and § 4) and the subsequent decay to the $2s$ state. The number density of atoms in the $2s$ state is

$$n_{2s} \simeq n_{1s}\Gamma_\beta/A_{2s1s}, \tag{4}$$

where $\Gamma_\beta = \int B_{3p1s}\phi_{13}(\nu)n_\alpha^{\rm src}(\nu)$ (in s^{-1}) is the transition rate to the $2s$ state owing to the Lyβ photons from the central source.

3. FINE-STRUCTURE LINES FROM THE EPOCH OF RECOMBINATION

The universe made a transition from fully ionized to nearly fully neutral at $z \simeq 1089$ (Spergel et al. 2007; Peebles 1968; Zeldovich et al. 1969). This transition is mainly accomplished by the two-photon decay of the $2s$ state and the slow redshifting of the Lyα photons which deplete the $2p$ state (see Peebles [1993] and references therein for a detailed discussion). The Saha ionization formula, valid for thermodynamic equilibrium conditions between the hydrogen level populations and free electrons, is a poor approximation for studying the epoch of recombination. A good approximation for studying this transition is to assume that all states, excepting $1s$, are in equilibrium with the CMBR (matter temperature to a very good approximation remains equal to the CMBR temperature throughout this transition; Seager et al. 1999; Peebles 1968). However, in this approximation where the matter temperature and all transitions, excepting the Lyα line, are in thermal equilibrium, the $2p$–$2s$ signal is unobservable, because the excitation temperature for this transition equals the background radiation temperature. Even though this second approximation might be useful for studying the evolution in ionization fraction, it will be strictly true only if the dominant mechanisms that determine the level populations of the $2p$ and the $2s$ states are either interaction with the CMBR photons or collisions between atoms. As there are a variety of other processes relevant to the determination of the level populations, for example, the free decay of either of the two states, a small deviation from equilibrium is expected in the $2s$ and $2p$ level populations, and it is our aim here to compute it.

One approach to this problem is to simultaneously solve for the level populations of the $2p$ and the $2s$ states as well as the change in the ionization fraction. However, assuming thermal equilibrium between these two states is a good approximation for solving the evolution of ionization. Therefore, the approach we have adopted is to solve for the evolution of ionization using the method of Peebles (1968) and use the resulting ionized/neutral fraction to solve for the populations of the $2s$ and the $2p$ states using detailed balance. The resulting equations are

$$\begin{aligned} f\alpha_B n_i^2 &+ cB_{2p2s}n_{2p}n_{\rm CMBR} + C_{\rm ps}n_i n_{2p} + A_{2p2s}n_{2p} \\ &+ n_{1s}A_{2s1s}\exp[-(B_1 - B_2)/(k_{\rm B}T_{\rm CMBR})] \\ &= A_{2s1s}n_{2s} + C_{sp}n_i n_{2s} + cB_{2s2p}n_{2s}n_{\rm CMBR} + \beta_c n_{2s} \end{aligned} \tag{5}$$

and

$$\begin{aligned} (1-2f)\alpha_B n_i^2 &+ cB_{2s2p}n_{2s}n_{\rm CMBR} + C_{\rm sp}n_i n_{2s} \\ &= A_{2p1s}n_{2p}/\tau_{\rm GP} + C_{ps}n_i n_{2p} + cB_{2p2s}n_{2p}n_{\rm CMBR} + \beta_c n_{2p}, \end{aligned} \tag{6}$$

where $B_1 = 13.6$ eV and $B_2 = 2.4$ eV are, respectively, the ionization potentials of the ground and the first excited states, $k_{\rm B}$ is the Boltzmann's constant, and $T_{\rm CMBR}$ is the temperature of the CMBR; β_c is the rate at which the CMBR photons cause a bound-free transition of electrons from the $n = 2$ state ($2s$ or $2p$ states). The various other terms in these equations have the same meanings as in the previous case. The main difference is that the CMBR photons and baryonic matter at the time of recombination are hot and dense enough to directly affect the level populations of the excited states by ionizing the excited state, and the two-photon capture to the excited state is not completely negligible (see Peebles [1968, 1993] for details of the different physical processes that are relevant at this epoch).

4. EXPECTED BRIGHTNESS OF THE FINE-STRUCTURE LINES

The brightness temperature in the $2p$–$2s$ transition is

$$\begin{aligned}\Delta T_b &\equiv T_b - T_{\rm CMBR} \\ &= g_\star \frac{n_{2p}(0) h_{\rm P} A_{2p2s} \lambda_e^2 (1+z)^2 c}{8\pi k_{\rm B} H(z)} \left(1 - \frac{T_{\rm CMBR}}{T_{\rm ex}}\right) \\ &\equiv \tau_{\rm ex}(T_{\rm ex} - T_{\rm CMBR}), \end{aligned} \quad (7)$$

where $n_{2p}(0) = n_{2p}/(1+z)^3$ and λ_e is the rest wavelength of the $2p$–$2s$ transition; $T_{\rm ex} = (h_{\rm P}\nu_e/k_{\rm B})[n_{2p}/(n_{2s} - n_{2p})]$ is the excitation temperature corresponding to the transition, where $h_{\rm P}$ is the Planck constant; $\tau_{\rm ex}$ is the optical depth of the source in the $2p$–$2s$ signal; $g_\star$ takes into account the selection rules for transitions between the $2p_{3/2}$ and $2s_{1/2}$ levels. Given the selection rules, there are three allowed transitions between these two states; they occur at frequencies $\nu \simeq 9852$, 9875, and 10029 MHz (see, e.g., Ershov 1987). The first two transitions are blended by the natural width of the line, which is approximately 100 MHz, but the third should be observable as a distinct line. This implies that $g_\star = 2/3$ if the observing frequency is $\simeq 9900/(1+z)$ MHz.

4.1. *Expectations for Signals from Cosmological* H II *Regions*

For computing the strength of this signal in a typical case, we adopt observed parameters of QSO SDSS J1030+0524 (see, e.g., Wyithe & Loeb 2004 and references therein), which is at redshift $z = 6.28$ and shows no detectable flux beyond the QSO Stromgren sphere; the radius of the Stromgren sphere has been estimated to be $R \simeq 4.5$ Mpc (Mesinger & Haiman [2004] argue that the size of the Stromgren sphere could be roughly 30% higher; this makes no essential difference to our results). It may be noted here that the QSO SDSS J1148+5251 also has similar parameters. We arrive at an estimate of $n_\alpha^{\rm src}(\nu)$ by assuming that L_α photons per second, with wavelengths corresponding to the Lyα transition, are emitted by the central source in an effective frequency range $\Delta\nu_{\rm source}$ and that the photons are absorbed at a radial distance R. This leads to

$$n_\alpha^{\rm src} = \frac{L_\alpha}{(4\pi R^2 c)} \frac{1}{\Delta\nu_{\rm src}}. \quad (8)$$

We assume typical values $L_\alpha = 10^{58}$ s^{-1}, $R = 4.5$ Mpc, and $\Delta\nu_{\rm src} = 30{,}000$ km s^{-1}. These lead to the following estimate for the transition rate, Γ_α, in the gas at the boundary of the Stromgren sphere arising due to the photons from the central QSO,

$$\Gamma_\alpha \simeq 8 \times 10^{-10}\ {\rm s}^{-1}. \quad (9)$$

It should be noted that the mean specific intensity of Lyα photons in the IGM would also give a nonzero signal. Assuming a mean specific intensity of $\simeq 10^{-21}$ ergs cm^{-2} s^{-1} Hz^{-1} sr^{-1} (this might be needed to couple the H I spin temperature to the matter temperature; see, e.g., Madau et al. 1997), the expected signal is many orders of magnitude smaller than we have computed from the outskirts of bright sources. If we assume that the central sources are continuum emitters in the frequency range between Lyα and Lyman-limit frequencies, we may assume similar parameter values for computing the expectations for Γ_β. In that case,

$$\Gamma_\beta \simeq \Gamma_\alpha \frac{f_\beta}{f_\alpha}, \quad (10)$$

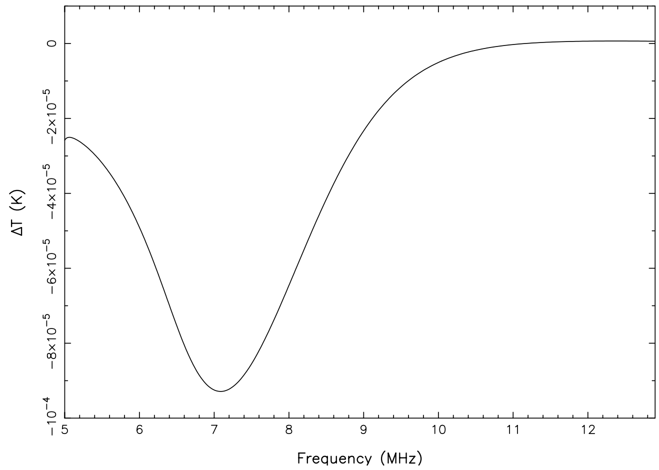

Fig. 1.—Expected brightness temperature decrement in the background radiation, owing to the fine-structure transition in gas at the recombination epoch, plotted vs. observing frequency.

where f_β and f_α are the oscillator strengths of the Lyβ and Lyα transitions, respectively. From equations (3), (4), and (10),

$$\frac{n_{2s}}{n_{2p}} \simeq \frac{f_\beta}{f_\alpha} \frac{A_{2p1s}}{A_{2s1s}\tau_{\rm GP}}. \quad (11)$$

For a completely neutral medium at $z \simeq 6.4$, $\tau_{\rm GP} \simeq 6 \times 10^5$, which may be used to show that $n_{2s} \gg n_{2p}$ for the parameters of IGM in the redshift range of interest, $6.4 \lesssim z \lesssim 10$. This implies that in the outskirts of the Stromgren sphere surrounding the QSO, the transition is expected to be observable as an absorption feature against the CMBR. Using equations (3) and (4) in equation (7), the observable brightness temperature is estimated to be

$$\Delta T_B \simeq -20\ \mu{\rm K} \left(\frac{f_{\rm neu}}{1}\right), \quad (12)$$

where $f_{\rm neu}$ is the neutral fraction of hydrogen outside the Stromgren sphere and might be close to unity.

We have adopted parameters typical of QSOs observed at the edge of the reionization epoch in estimating the above temperature decrement. The main uncertainty above is in the estimation of the "Lyα" flux from the central source, and as the observed temperature decrement is directly proportional to this flux from the central source, this constitutes a major uncertainty in reliably computing the expectations for the signal. For QSOs that have strong Lyα and Lyβ lines, the blueward side of the lines will be strongly absorbed in the medium just beyond the Stromgren sphere (and this has been observed to happen in many cases), provided that the blueward side photons have not been redshifted to frequencies smaller than the Lyα frequency while traversing the Stromgren sphere. In the case of QSOs that have large line fluxes and small Stromgren spheres, the Lyα luminosity L_α might be underestimated.

4.2. *Expectations for the Fine-Structure Line from Recombination*

Using equations (5) and (6), the level populations of the $2p$ and the $2s$ states may be computed. Solving for the level populations and using equation (7), we have computed the expected signal from the recombination epoch; the expected signal is shown in Figure 1. The fine-structure line transition is expected to be an absorption feature, with a maximum temperature decrement of order 100 μK at an observing frequency of about 7 MHz. The width of the decrement in frequency space corresponds to a redshift span $\Delta z \simeq 200$, which is roughly the width of the visibility function at recombination.

5. PROSPECTS FOR THE DETECTION OF THE COSMOLOGICAL FINE-STRUCTURE LINES

As equation (12) shows, a detection of the fine-structure line in the outer regions of cosmological H II regions is potentially a probe of the cosmological neutral hydrogen density in the vicinity of QSOs in the reionization epoch and might be a tool for the investigation of the evolution in the neutral fraction with cosmic epoch through the reionization era. Given the cosmological importance of such a measurement and the fact that there does not exist many reliable methods for the detection of H I at high redshifts (see, e.g., Barkana & Loeb 2001), the detection of the $2p$–$2s$ line transition in the cosmological context assumes additional importance.

In deriving equation (7), the Hubble expansion was assumed to be the only cause for the velocity width in the observed line. However, an important contribution to the velocity dispersion in the line in this case is the natural width of the fine-structure line; owing to the rapid decay of the $2p$ state, the natural Lorentzian width in the rest frame of the gas is $V_{\rm Lor} \simeq 100$ MHz (see, e.g., Dennison et al. 2005). Therefore, the peak brightness in the observed line profile might be suppressed by a factor $V_{\rm exp}/V_{\rm Lor}$, where $V_{\rm exp}$ is the line-of-sight peculiar velocity dispersion owing to the Hubble flow across the region being observed. However, in the case $V_{\rm exp} \gtrsim V_{\rm Lor}$ this suppression has a negligible effect. For QSOs at redshift $z \approx 6.5$, the natural Lorentzian width of the fine-structure line is equivalent to the peculiar Hubble flow across a proper line-of-sight distance of $\simeq$4 Mpc. This distance is approximately the size of the Stromgren spheres around QSOs at that redshift; therefore, the natural width of the line does not significantly diminish the expectations, given by equation (7), for the peak brightness temperature. A second inference is that in the case of a QSO that has a smaller Stromgren sphere, the increase in mean brightness temperature is roughly proportional to the inverse of the radius of the Stromgren sphere, $1/R$, and not $1/R^2$.

Another assumption that was made while deriving equation (7) is that the only radio frequency radiation that needs to be considered for the determination of the level populations and brightness temperature decrement is the CMBR. It may be that the central ionizing source, which may be a QSO, is radio loud. There may also exist radio sources behind the observed Stromgren sphere and within the angular region over which the Lyα flux from the QSO is appreciable. Equation (7) may be modified, to account for this, by replacing $T_{\rm CMBR}$ with $T_{\rm CMBR} + T_B$, where T_B is the brightness temperature of the radio source at wavelengths corresponding to the rest frequency of the transition, which is $\simeq$9 GHz. The result would be to enhance the brightness temperature of the line signal by $\simeq T_B/T_{\rm CMBR}$. If the radio source is unresolved, it is appropriate to estimate the expected signal in terms of the optical depth. The optical depth corresponding to equation (12) is $\simeq 10^{-6}$.

We now discuss the feasibility of the detection of the fine-structure line toward SDSS J1030+0524. We shall assume that the neutral fraction, $f_{\rm neu}$, outside the Stromgren sphere of this QSO is unity, consistent with the measured GP trough. The redshifted fine-structure line would be expected at $\simeq$1.36 GHz. The observed line "width," considering natural broadening and the Hubble flow across the Stromgren sphere, is expected to be approximately $\Delta\nu \simeq 100$ MHz$/(1 + z)$; using $z = 6.28$, we obtain $\Delta\nu = 13.7$ MHz. The fine-structure line would be expected to originate in a shell that is roughly the size of the Stromgren sphere and fall off as $1/r^2$ beyond this shell, where r is the distance from the QSO. In the case of SDSS J1030+0524, the angular size of the Stromgren sphere is expected to be $15'$.

The frequency of the redshifted line is in the observing bands, and the line width is within the spectral line capabilities, of several currently operational telescopes. However, large-collecting-area arrays like the Giant Metrewave Radio Telescope (GMRT) have large-aperture antennas of 45 m diameter and, therefore, poor surface brightness sensitivity for such extended structures. The Australia Telescope Compact Array (ATCA), with 22 m antennas, has reasonable brightness sensitivity for this problem and, additionally, operates with 128 MHz bandwidths in spectral-line mode. At $\nu \simeq 1.3$ GHz, the ATCA has a system temperature $T_{\rm sys} \simeq 25$ K and antenna efficiency $K = 0.1$ K Jy^{-1}. The five movable 22 m diameter antennas may be configured into an ultracompact two-dimensional close-packed H75 (75 m maximum baseline) array, and this would yield a number of baselines sensitive to the $15'$ scale fine-structure line signal. The brightness sensitivity of this array, for a $8'$ scale structure, is 400 μK in 6 hr integration time. The brightness sensitivity in Fourier synthesis images could be enhanced somewhat, by factors of a few, by appropriately weighting the baselines to match the synthetic beam to the expected structure scale; however, the required integration times for a detection are still in the ballpark of 10^3 hr.

The detection of the signal from cosmological H II regions, however, might be feasible using facilities under construction or planned for the near future, like the Extended New Technology Demonstrator (xNTD) in Australia or the Square Kilometer Array (SKA). These arrays would have smaller antenna sizes, making the detection of these large-angular-scale structures detectable in interferometers, and significantly greater numbers of antennas, giving greater numbers of short baselines that would usefully respond to the large-angular-scale fine-structure line. However, the array configuration designs would need to factor in the extraordinary brightness temperature sensitivity requirements for this demanding observation.

As an example, the SKA might consist of about 10^4 12–15 m class reflector antennas, with aperture efficiency of 60%, and the system temperature at 1.4 GHz might be about 20 K. Assuming that the visibilities are optimally weighted so that the synthetic beam of the Fourier synthesis array is matched to the source size of about $15'$ FWHM, the line strength would have a peak of about 20 μJy. A 5 σ detection of the fine-structure line toward a typical QSO, in a reasonable integration time of about 1 hr, will require that about 250 baselines (just 0.001% of the total) be within about 50 m.

The signal from the recombination epoch is observable in the frequency band 6–8 MHz as a broad decrement in the brightness temperature of the extragalactic background sky and would be extremely small compared to the orders of magnitude more intense Galactic nonthermal emission as well as the average low-frequency background brightness temperature arising from the numerous extragalactic radio sources. This decrement may be considered to be a distortion to the CMBR spectrum at long wavelengths and would be an all-sky cosmological signal. However, owing to the ionosphere, the frequency range in which this feature is expected to appear is too low a radio frequency to be easily accessible using ground-based observatories. Therefore, even though the observation of this signal might be yet another tool for probing the epoch of recombination, new custom-made instruments, which should presumably operate from space and above the ionosphere, will need to be built if a detection of this signal is to be attempted. An additional cause for concern is that the Galactic and extragalactic background radiations might have low-frequency spectral turnovers at these frequencies, as a result of free-free absorption as well as synchrotron self-absorption, and these would result in significant spectral features in the band that would require a careful modeling in order to detect any CMBR decrement feature arising from fine-structure transition absorption. Interference from terrestrial man-made transmitters, as well as from auroral phenomena and solar system objects would also be an issue.

6. SUMMARY AND DISCUSSION

We have discussed the possibility of detecting the hitherto undetected fine-structure line of $2p$–$2s$ transition in two cosmological settings: the epoch of recombination at $z \simeq 1089$ and the epoch of reionization at $z \simeq 6$–15.

The expectations for the line from the environments of ionization sources in the epoch of reionization are interesting and worthy of attention as a novel tool for the investigation of the reionization process and the cosmological evolution of the gas. The signal is expected to be observable as an extended and weakly absorbing source, which causes a decrement in the brightness of the background CMBR, and may be detected by interferometers as a negative source akin to the Sunyaev-Zeldovich decrements observed along the lines of sight through hot gas in clusters of galaxies. Detection of the $2p$–$2s$ line signal from the outskirts of cosmological H II regions at different redshifts within the reionization era may serve to determine the neutral fraction of the medium during the epoch of reionization, which is a quantity of significant interest in modern cosmology. In particular, we have computed a representative signal strength by adopting parameters typical of a QSO observed at $z \simeq 6$. These QSOs show GP troughs in their spectra; however, owing to the weakness of the GP test, the spectra have only been useful in setting weak limits on the neutral fraction (constraining the neutral fraction to be $\gtrsim 10^{-3}$) outside the observed H II region. We have shown that for a fully neutral medium, the line peak may reach $\simeq 20\ \mu$K, which is potentially observable by radio interferometers that are being designed today. Other interesting probes of the reionization epoch include detecting O I line from this epoch (Hernandez-Monteagudo et al. 2006).

It is of interest to compare the relative difficulty associated with detecting the neutral hydrogen in this indirect way, using the redshifted fine-structure line, with direct imaging of the redshifted 21 cm line from neutral hydrogen during the epoch of reionization (see, e.g., Sethi 2005; Zaldarriaga et al. 2004). The "all-sky" H I signal might be detectable with a peak strength of $\simeq 50$ mK; however, it could be very difficult to detect owing to calibration and foreground contamination issues (e.g., Zaldarriaga et al. 2004; Shaver et al. 1999, and references therein). A better approach might be to attempt to detect the fluctuating component of the sky signal, which could have peak intensities of approximately a few mK at observing frequencies of $\nu \simeq 100$–200 MHz (Zaldarriaga et al. 2004; Shaver et al. 1999). This translates to roughly the same signal strength (specific intensity) as we have obtained in equation (12) for the fine-structure line. This is not entirely unexpected; even though the level populations of the excited states are much smaller as compared to the ground level population needed in computing the H I signal, the A-coefficient of the fine-structure transition we consider here is roughly 8 orders of magnitude larger than the H I hyperfine-transition A-coefficient. Currently, there are many ongoing and planned radio interferometer experiments for the detection of the redshifted H I emission/absorption from the epoch of reionization (e.g., Pen et al. 2004; the LOFAR project[1]; the MWA project[2]). The detection of the H I signal, first, is more difficult, requiring greater sensitivity, because the typical frequency width of the signal $\simeq 0.5$ MHz (Zaldarriaga et al. 2004), which is far smaller than the typical width expected for the fine-structure line ($\simeq 10$ MHz). Second, the fine-structure line is expected at 1.4 GHz, where the sky background temperatures are significantly lower than in the 100–200 MHz band, making the telescope system temperatures lower. An additional advantage of the indirect detection is that the redshifted fine-structure line appears at higher frequencies ($\gtrsim 1$ GHz), which are relativity free of interference as compared to the low-frequency band of 100–200 MHz. The main advantage of the direct detection is that unlike the method we suggest here, it is independent of the existence of strong Lyα emitters at high redshifts. Another advantage of direct detection is that it may be detected "statistically," and such a detection might be achieved with greater ease than direct "imaging" (e.g., Zaldarriaga et al. 2004; Bharadwaj & Sethi 2001); however, the foreground subtraction problem becomes a severe constraint for a statistical detection. To summarize, if strong Lyα-emitting sources are present at high redshifts, they would facilitate the indirect detection of the neutral hydrogen via enabling the detectability of the fine-structure line. The imaging issues and problems associated with the detection of this signal appears to be less of a challenge as comparison to the direct detection of redshifted H I from those epochs.

Another astrophysical context in which the fine-structure line might be detectable is the environments of high-redshift galaxies, which are strong Lyα emitters. As equation (12) shows, the observed signal depends on the neutral fraction outside the Stromgren sphere. Therefore, detection of this signal would constitute an alternate probe of the neutral fraction of the IGM at large redshifts.

The aim of this work has been to examine the detectability of the fine-structure line in cosmological contexts, to point out the cosmological significance of detections, and spawn work that may refine the modeling presented herein and improve the case for appropriate design of future telescopes, which might enable the detection of the fine-structure line toward multiple sources in the reionization era.

One of us (S. K. S.) would like to thank Jens Chluba for many useful discussions and to Zoltan Haiman for many useful comments on the manuscript.

[1] See Web site at http://www.lofar.org.
[2] See Web site at http://www.haystack.mit.edu/ast/arrays/mwa/site/index.html.

APPENDIX

PHOTON DISTRIBUTION FUNCTION

The evolution of the photon distribution function, neglecting the effect of scattering off moving atoms, is (see, e.g., Rybicki & Dell'Antonio 1994)

$$\frac{\partial n_\nu}{\partial t} - \nu H \frac{\partial n_\nu}{\partial \nu} = A_{2p1s} n_{2p} \phi_\nu - c B_\nu n_{1s} \phi_\nu n_\nu, \quad (13)$$

where $B_\nu = 3/(8\pi)c^2/\nu_\alpha^2 A_{2p1s}$ and $H = \dot{a}/a$. The equation may be written as

$$\frac{\partial n_\nu}{\partial t} - \nu H \frac{\partial n_\nu}{\partial \nu} = -1/\tau(n_\nu - n_\star), \tag{14}$$

where $\tau = 1/(cB_\nu\phi_\nu n_{1s})$ and $f_\star = A_{2p1s}n_{2p}/(B_\nu n_{1s}c)$. The equation above lends itself to a ready interpretation. If the second term on the left-hand side (which is owing to the expansion of the universe) was absent, the distribution function will approach $f_\star$ on a timescale $\simeq\tau$, where $\tau \simeq 3n_{1s}$ s $\ll \dot{a}/a$ (n_{1s} here has units cm^{-3}). It may be readily verified that for the recombination epoch and also the epoch of reionization, $\tau \ll 1/H$, excepting when the neutral fraction of the medium is very small. We work here with the assumption that the neutral fraction is always large enough so that $\tau \ll 1/H$. In this case the solution to equation (14) may be simplified; to leading order the distribution function approaches $f_\star$ and the first-order term (of the order of τH) represents the slow time variation of the distribution function owing to the expansion of the universe. In this approximation, one may write the solution to equation (14) as

$$n_\nu \simeq n_\star - \tau H(t) n_\star. \tag{15}$$

Using these equations, we may proceed to prove the contention that the net effect of the "trapping" of Lyα photons is to reduce the decay time of the $2p$ state by a factor $\tau_{\rm GP}$. The $1s$–$2p$ transition rate, which is given by equation (6), may be solved using equation (15),

$$cn_{1s}\int B_\nu\phi(\nu)n(\nu)d\nu \simeq A_{2p1s}n_{2p} - A_{2p1s}n_{2p}/\tau_{\rm GP}. \tag{16}$$

The first term on the right-hand side cancels with the decay term of the $2p$ state on the right-hand side of equation (6), and therefore, the net effect of the scattering of recombination photons is to reduce the decay time of the $2p$ state by a factor $\tau_{\rm GP}$. It may be pointed out that the condition needed to derive the above expression roughly translates to the condition that $\tau_{\rm GP} \gg 1$. For the reionization case, this requires that the neutral fraction $\gtrsim 10^{-5}$. In the case of primordial recombination, it leads to an even weaker condition that the neutral fraction is $\gtrsim 10^{-7}$.

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