EVLATION O THE IONIZING BACKGROUND AND THE EPOCH OF REIONIZATION FROM THE SPECTRA OF z \sim 6 QUASARS

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ABSTRACT

We study the process of cosmic reionization and estimate the ionizing background in the intergalactic medium (IGM) using the Lyman series absorption in the spectra of the four quasars at 5.7 < z < 6.3 discovered by the Sloan Digital Sky Survey. We derive the redshift evolution of the ionizing background at high redshifts, using both semianalytic techniques and cosmological simulations to model the density fluctuations in the IGM. The existence of the complete Ly\alpha Gunn-Peterson (GP) trough in the spectrum of the z = 6.28 quasar SDSS 1030+0524 indicates a photoionization rate (\Gamma_{-12} in units of 10^{-12} s^{-1}) at z \sim 6 lower than 0.08, at least a factor of 6 smaller than the value at z \sim 3. The Ly\beta and Ly\gamma GP troughs give an even stronger limit \Gamma_{-12} \lesssim 0.02 due to their smaller oscillator strengths, indicating that the ionizing background in the IGM at z \sim 6 is more than 20 times lower than that at z \sim 3. Meanwhile, the volume-averaged neutral hydrogen fraction increases from 10^{-5} at z \sim 3 to greater than 10^{-3} at z \sim 6. At this redshift, the mass-averaged neutral hydrogen fraction is larger than 1\%; the mildly overdense regions (\delta > 3) are still mostly neutral, and the comoving mean free path of ionizing photons is shorter than 8 Mpc. Comparison with simulations of cosmological reionization shows that the observed properties of the IGM at z \sim 6 are typical of those in the era at the end of the overlap stage of reionization when the individual H II regions merge. Thus z \sim 6 marks the end of the reionization epoch. The redshift of reionization constrains the small-scale power of the mass-density fluctuations and the star-forming efficiency of the first generation of objects.

Key words: intergalactic medium — quasars: absorption lines

1. INTRODUCTION

After the recombination epoch at z \sim 1500 the universe remained mostly neutral until the first generation of stars and quasars ionized the intergalactic medium (IGM) and ended the cosmic “dark ages” (e.g., Rees 1998). Popular cosmological models predict this reionization to have occurred at redshift between 6 and 20 (e.g., Gnedin & Ostriker 1997; Chiu & Ostriker 2000; Gnedin 2000; Ciradi et al. 2001; Razoumov et al. 2001, and references therein). When and how the universe reionized is one of the fundamental questions of modern cosmology (for a complete review, see Barkana & Loeb 2001 and Loeb & Barkana 2001). Observations of absorption spectra of luminous sources at high redshifts approaching the reionization epoch provide the best observational probes of reionization to date.

Recent discoveries of luminous quasars at z > 5.7 (Fan et al. 2000, 2001c), using imaging data from the Sloan Digital Sky Survey (SDSS, York et al. 2000), enable us to study the state of the IGM at redshift up to 6. Observations of Fan et al. (2001c, hereafter Paper I), Becker et al. (2001, hereafter Paper II), and Pentericci et al. (2002, hereafter Paper III) show that the Ly\alpha absorption due to neutral hydrogen in the IGM increases dramatically toward high redshifts. In particular, Keck and VLT spectroscopy of high-redshift quasars (in Papers II and III, respectively) show the first observation of a complete Gunn-Peterson (GP) trough (Shklovsky 1964; Scheuer 1965; Gunn & Peterson 1965) in the spectrum of the z = 6.28 quasar SDSSp J103027.10+052455.0 (SDSS 1030+0524 for brevity, where no flux is detected over a 300 \AA region immediately blueward of the Ly\alpha emission line. The flux decrement between the red and blue sides of the Ly\alpha emission line is larger than 150, corresponding to an effective optical depth to Ly\alpha photons (\tau_{\rm eff}) larger than 5 at z_{\rm abs} \sim 6.05. The existence of a Ly\beta GP trough in the Keck spectrum imposes an even stronger limit on the effective equivalent Ly\alpha optical depth (\tau_{\rm eff} > 20, Paper II). Djorgovski et al. (2001) also observed a particularly dark region of length \sim 5 Mpc at z \sim 5.4, along the line of sight to the z \sim 5.8 quasar SDSS 1044–0125.

This is a theoretical companion paper to Papers I, II, and III. In this paper we use the absorption measurements presented in the previous papers to calculate the evolution of the ionizing background and the neutral hydrogen fraction in the IGM and constrain the epoch of reionization. The outline of this paper is as follows. In § 2 we summarize the observed evolution of the average Ly\alpha absorption along the lines of the sight to the sample of quasars in Paper I, based on the spectra presented in Papers II and III, and discuss the errors on these measurements. In § 3 we calculate the ionizing background at different redshifts using the average Ly\alpha absorption, taking into account the inhomogeneity of the IGM using a model for the overdensity distribution. We derive stronger constraints on the ionizing background in the GP trough region of SDSS 1030+0524 at z_{\rm abs} \sim 6 using

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the Ly$\beta$ and Ly$\gamma$ absorption troughs. In § 4 we use N-body simulations to model the high-redshift Ly$\alpha$ forest spectra and compare the statistical properties of the simulated and observed spectra. In § 5 we calculate the redshift evolution of the volume and mass-weighted neutral hydrogen fractions in the IGM, and show how the reionization penetrates into progressively overdense regions at lower redshifts as the mean free path of ionizing photons increases. These calculations are used to constrain the epoch of reionization. We discuss the prediction of structure formation models for the reionization epoch in § 6 and outline future observations needed to probe deep into and beyond the reionization epoch.

Throughout the paper, we adopt a $\Lambda$-dominated cold dark matter (LCDM) cosmology, with $\Omega_m = 0.35$, $\Lambda = 0.65$, $h = 0.65$, and $\Omega_b h^2 = 0.02$, unless otherwise noted.

2. EVOLUTION OF NEUTRAL HYDROGEN ABSORPTION

The Gunn & Peterson (1965) optical depth to Ly$\alpha$ photons is

$$\tau_{GP} = \frac{\pi e^2}{m_e c} f\lambda H^{-1}(z)n_{H_1},$$

where $f\lambda$ is the oscillator strength of the Ly$\alpha$ transition, $\lambda = 1216$ Å, $H(z)$ is the Hubble constant at redshift $z$, and $n_{H_1}$ is the density of neutral hydrogen in the IGM. At high redshifts $H(z) \approx h\Omega_m^{1/2}/(1 + z)^{3/2}$, and the GP optical depth for a uniformly distributed IGM can be rewritten as

$$\tau_{GP}(z) = 1.8 \times 10^5 h^{-1} \Omega_m^{-1/2} \left[ \frac{\Omega_b h^2}{0.02} \right] \left[ \frac{1 + z}{7} \right]^{3/2} \left( \frac{n_{H_1}}{n_{H_1}} \right).$$

In reality, the IGM is highly inhomogeneous, and regions with different densities will be in different ionization states and have different GP optical depths (§ 3). The IGM density distribution needs to be taken into account when estimating the averaged neutral fraction from the observations (§ 5).

Figure 1 shows the observed evolution of the average Ly$\alpha$ absorption in the high-redshift quasar spectra as a function of redshift, both in terms of the transmitted flux ratio $\mathcal{F}$, defined as the ratio of observed and unabsorbed continuum fluxes (see Paper I),

$$\mathcal{F}(z_{abs}) \equiv \left\langle \frac{f_{\nu}^{obs}}{f_{\nu}^{con}} \right\rangle,$$

and the effective GP optical depth, defined as

$$\tau_{GP}^{eff} \equiv -\ln(\mathcal{F}).$$

The measurements shown in Figure 1 at $z_{abs} < 4.5$ are taken from McDonald & Miralda-Escudé (2001). The measurements at $4.8 < z_{abs} < 5.6$ are based on the spectra of Paper II and III. In this figure we have averaged the measurements along four lines of sight into four different redshift bins, $z_{abs} = (4.8–5.2), (5.2–5.6), (5.6–5.95)$, and $(5.95–6.15)$, respectively, and include the systematic errors due to the extrapolation of the intrinsic quasar continuum $\sigma$(con). In Papers I, II, and III, we assumed a power-law continuum $f_{\nu} \propto \nu^{-\alpha}$, where $\alpha = 0.5$, normalized at rest-frame wavelength 1280 Å. Therefore,

$$\sigma_{con} = \left| \log \left[ \frac{1216(1 + z_{abs})}{1280(1 + z_{em})} \right] \right| \sigma(\alpha) \times \mathcal{F}.$$ (5)

For $\sigma(\alpha) \sim 0.4$ (Fan et al. 2001a) the typical relative error is $\sigma_{con}/\mathcal{F} \sim 0.05$. (3) Sample variance $\sigma$(sample). This quantity can be roughly estimated using a model in which the Ly$\alpha$ clouds are distributed randomly in the IGM (Zuo 1992). We calculate $\sigma$(sample) using equations (8) and (9) of [Paper].

![Fig. 1.—Evolution of Ly$\alpha$ absorption based on the observations of four quasars at $z > 5.7$ in Fan et al. (2001c), Becker et al. (2001), and Paper III. The results at $z_{abs} < 5.6$ are averaged over four lines of sight, and the error bars include contributions from photon noise, uncertainty in the intrinsic quasar continuum, and an estimate of the sample variance.](image-url)
Zuo (1992), assuming that the column-density distribution of the Lyα clouds follows a power law \( N(N_{\text{HI}}) \propto N_{\text{HI}}^{\beta} \), with \( \beta = -1.5 \), and the Doppler width of the Lyα forest line is \( b = 30 \, \text{km} \, \text{s}^{-1} \). In a redshift window \( \delta z = 0.4 \), \( \sigma(\text{sample}) \sim 0.025 \) for \( \mathcal{F} \sim 0.10 \), while for \( \delta z = 0.2 \), \( \sigma(\text{sample}) \sim 0.002 \) for \( \mathcal{F} \sim 0.006 \). Therefore, for all the measurements, except at the highest redshift (\( z_{\text{abs}} \sim 6 \)), the error term is dominated by sample variance. This conclusion is consistent with the scatter of \( \mathcal{F} \) measurements from different lines of sight at the same redshift, which is larger than that expected from photon noise alone. For \( \mathcal{F} \ll 1 \), the contribution from photon noise becomes important. The error bars in Figure 1 represent the sum of the three error terms, added in quadrature.

3. EVOLUTION OF THE IONIZING BACKGROUND

The Lyα forest arises from absorption of Lyα photons by neutral hydrogen gas in the intergalactic medium (eq. [2]). Assuming local photoionization-recombination equilibrium, we have

\[
n_{\text{H}i} n_e \Gamma = n_{\text{H}i} n_e \alpha(T),
\]

where \( n_{\text{H}i}, n_{\text{H}e} \) and \( n_e \) are the local densities of neutral, ionized hydrogen and electrons in the IGM, respectively, \( \Gamma \) is the photoionization rate, and \( \alpha(T) \) is the recombination coefficient at temperature \( T \) (Abel et al. 1997),

\[
\alpha(T) = 4.2 \times 10^{-13} (T/10^4 \text{K})^{-0.7} \, \text{cm}^3 \, \text{s}^{-1}.
\]

The photoionization rate \( \Gamma \) is related to the ionizing background flux \( J_\nu \) by

\[
\Gamma = 4\pi \int \frac{J_\nu}{h\nu} \sigma_\nu \, d\nu,
\]

where the integral includes the ionizing photons from the H I Lyman limit to the He II Lyman limit, assuming that the He in the IGM is singly ionized and \( \sigma_\nu \) is the H I cross-section of ionizing photons, \( \sigma_\nu \propto \nu^{-3} \). For a power-law spectrum of the ionizing background dominated by quasars, \( J_\nu \propto \nu^{-1.8} \) (Madau, Haardt, & Rees 1999), we find \( \Gamma_{-12} = 2.5 J_{-21} \), where \( \Gamma_{-12} \) is the photoionization rate in units of \( 10^{-12} \, \text{s}^{-1} \) and \( J_{-21} \) is the ionizing flux at the Lyman limit in units of \( 10^{-21} \text{ergs} \, \text{cm}^{-2} \, \text{s}^{-1} \, \text{Hz}^{-1} \, \text{sr}^{-1} \). If the ionizing background at high redshifts is dominated by massive stars (e.g., Paper I) with \( J_\nu \propto \nu^{-3} \) ( Barkana & Loeb 2001), we find \( \Gamma_{-12} = 1.5 J_{-21} \).

If the IGM is mostly ionized by a uniform ionizing background, the evolution of the optical depth can be expressed as (Weinberg et al. 1997)

\[
\tau_{\text{GP}} \propto \frac{(1+z)^6 (\Omega_b h^2)^2 \alpha(T)}{\Gamma(z)} \left[ 1 + \frac{0.05}{\Gamma_{-12}(z)} \right] \Delta^2.
\]

The Lyα absorption increases rapidly with increasing redshift even if the ionizing background remains constant with redshift. McDonald & Miralda-Escudé (2001) estimate the ionizing background at \( z < 5.2 \) by comparing the observed transmitted flux ratio with that of artificial Lyα forest spectra created from cosmological simulations. In this section we develop a semianalytic model, which uses an empirical model for the density distribution derived from hydrodynamical simulations similar to that of McDonald & Miralda-Escudé (2001), and use it to estimate the ionizing background from the Lyα observations (§3.1). Then we use the Lyα and Lyγ forest GP trough measurements in the spectrum of SDSS 1030+0524 to put stronger constraints on the background at \( z \sim 6 \) (§3.2).

3.1. Semianalytic Modeling of Lyα Absorption

The Lyα forest arises from low-density gas in the IGM that is in approximate thermal equilibrium between photoionization heating by the UV background and adiabatic cooling due to Hubble expansion (see, e.g., Bi 1993; Cen et al. 1994; Zhang, Anninos, & Norman 1995; Hernquist et al. 1996; Hui & Gnedin 1997). The neutral hydrogen fraction and therefore the GP optical depth depend on the local density of the IGM. We define the fractional density of the IGM as \( \Delta \equiv \rho/\langle \rho \rangle \) (\( \delta + 1 \), where \( \delta \), the density contrast, is the departure of the local density from the mean density, in units of the mean density). For a region of IGM with density \( \Delta \)

\[
\tau(\Delta) \propto \frac{(1+z)^4 \rho^{0.2} \alpha(T) [\Gamma(\Delta, z) \Omega_m^{0.5}]}{h^4 (\gamma + 1)} \Delta^2.
\]

The dependence on \( \Delta^2 \) arises because \( \tau \propto n_{\text{H}i} \), which is proportional to \( n_{\text{H}i}^{\beta} \) (eq. [6]) and proportional to \( \Delta^2 \) for a highly ionized IGM. The temperature of the IGM is determined by photoionization-recombination equilibrium, which leads to a power-law relation between temperature and density of the form \( T = T_0 \Delta^\gamma \), with \( T_0 \sim 1 - 2 \times 10^4 \text{K} \), and \( \gamma \sim 0.1 \) (e.g., Hui & Gnedin 1997). Following McDonald & Miralda-Escudé (2001), we assume an uniform ionizing background and \( \gamma = 0 \). Thus

\[
\tau_0(\Delta) = \frac{(1+z)^{4.5} \rho^{0.2} \alpha(T) [\Gamma_{-12}(z)]}{h^4 (\gamma + 1)} \Delta^2.
\]

We determine \( \tau_0 \). The mean transmitted flux ratio \( \mathcal{F} \) can be calculated as

\[
\mathcal{F} = \left< e^{-\tau} \right> = \int_0^\infty e^{-\tau} p(\tau) d(\tau) ,
\]

where \( p(\tau) \) is the distribution function of the density of the IGM. For an inhomogeneous IGM, using the definition of effective optical depth in equations (4) and (12), we have \( \mathcal{F} = e^{-\tau_{\text{eff}}} = \left< e^{-\tau} \right> > e^{-\tau_{\text{eff}}} \), thus \( \tau_{\text{eff}} < \left< \tau \right> \).

We calculate \( \mathcal{F} \) using a parametric form for the volume-weighted density-distribution function \( p(\Delta) \) (Miralda-Escudé, Haehnelt, & Rees 2000):

\[
p(\Delta) = A \exp \left[ -\frac{(\Delta^{-2/3} - C_0)^2}{2(2\delta_0/3)^2} \right] \Delta^{-\beta}.
\]

where \( \delta_0 = 7.61/(1+z) \), and \( \beta, C_0 \), and \( A \) are numerical constants given in Table I of Miralda-Escudé et al. (2000) at several redshifts. This distribution is derived assuming that the initial density fluctuations form a Gaussian random field, the gas in voids is expanding at a constant velocity, and the density field is smoothed on the Jeans scale of the photoionized gas. It reproduces well the density distribution of the photoionized gas in the LCDM hydrodynamic simulations of Miralda-Escudé et al. (1996). In the following
calculations we linearly interpolate \( \beta \) in redshift and set the constants \( C_0 \) and \( A \) so that both the total volume and mass are normalized to unity. Finally, we find the numerical constant in equation (11) to be \( \tau_0 = 82 \) by requiring the resulting \( \Gamma \) to match the measurement of McDonald & Miralda-Escudé (2001) at \( z = 4.5 \). Note that the quantity we are really constraining with Ly\( \alpha \) absorption is the combination \( \Gamma \Omega_m^2 \Omega_h^2 \sigma_0^2 T_{\text{c}}^{-2} \) (eqs. [10] and [11]). There are still considerable uncertainties in the values of \( \Omega_m, \Omega_h, h, \) and \( T_0 \), all of which translate to uncertainties in the normalization of \( \Gamma_{-12} \). However, in the era after reionization, the temperature of the IGM (\( T_0 \)) is not a strong function of redshift. Therefore, the redshift evolution rate of \( \Gamma_{-12} \) derived here is robust to uncertainties in the normalization. Note that McDonald & Miralda-Escudé (2001) assume that \( T_0 = 2 \times 10^4 \) K, \( \Omega_b h^2 = 0.02, h = 0.65, \) and \( \Omega_m = 0.4 \). The \( \Gamma_{-12}(z) \) would be ~40\% higher for \( T_0 = 10^4 \) K (e.g., Schaye et al. 2000). McDonald & Miralda-Escudé (2001) show that the estimate of \( \Gamma_{-12} \) is quite insensitive to the slope of the \( \Delta-T \) relation (i.e., the value of \( \gamma \)).

Also note that the effective optical depth here is calculated by integrating in real space rather than in redshift space. We compared the results from \( N \)-body simulations using redshift and real space coordinates (\( \S \) 4) and found that the effect of ignoring peculiar velocities on the resulting optical depth is negligible.

Figure 2 shows the evolution of the photoionization rate \( \Gamma_{-12} \) based on the observations presented in Figure 1, including the measurements at \( z < 4.5 \) from McDonald & Miralda-Escudé (2001). The ionizing background declines toward higher redshift, from \( \Gamma_{-12} \sim 0.5 \) at \( z_{\text{abs}} \sim 3 \), to \( \Gamma_{-12} \sim 0.2 \) at \( z_{\text{abs}} \sim 5 \), and to \( \Gamma_{-12} \sim 0.12 \) at \( z_{\text{abs}} \sim 5.8 \). Note that \( \Gamma \) remains roughly constant between \( z \sim 4.4-5.5 \), albeit with rather large error bars on individual measure-

3.2. Constraints from Ly\( \beta \) and Ly\( \gamma \) Troughs in SDSS 1030+0524

As pointed out in Paper II, because of its much smaller oscillator strength and the fact that the flux scales exponentially with oscillator strength, one can put a stronger limit on the neutral hydrogen fraction in the GP trough from the Ly\( \beta \) transition than from the Ly\( \alpha \) transition, even after allowing for the foreground Ly\( \alpha \) absorption at lower redshifts. In this subsection, we put a much stronger limit on \( \Gamma_{-12} \), based on the Ly\( \beta \) measurement of Paper II. We also attempt to use the Ly\( \gamma \) trough to set a limit on \( \Gamma_{-12} \), although it suffers from larger uncertainties in the modeling due to the presence of both foreground Ly\( \beta \) and Ly\( \alpha \) absorption, and it is contaminated by Ly\( \delta \) absorption.

The Ly\( \beta \) absorption at redshift \( z_\beta \) overlaps with the foreground Ly\( \alpha \) absorption at the redshift

\[
\frac{z_\alpha}{z_\beta} = \left( \frac{\lambda_\alpha/\lambda_\beta}{1+\beta_0} \right) (1 + \beta_0) - 1 ,
\]

where \( \lambda_{\alpha,\beta} \) are the rest-frame wavelengths of the Ly\( \alpha \) and Ly\( \beta \) transitions, respectively. The observed total optical depth at the corresponding wavelength is \( \tau_{\text{total}} = \tau_{\alpha}(z_\alpha) + \tau_{\beta}(z_\beta) \). For a fractional density \( \Delta \),

\[
\tau_{\beta}(\Delta) = \frac{f_\beta}{f_\alpha} \tau_\alpha(\Delta) = \tau_\alpha(\Delta)/5.27 ,
\]

where \( f_\alpha \) and \( f_\beta \) are the oscillator strengths of Ly\( \alpha \) and Ly\( \beta \) transitions, respectively, and \( \tau_\alpha(\Delta) \) can be calculated using equation (11). Therefore, the observed transmitted flux ratio in the presence of both Ly\( \alpha \) and Ly\( \beta \) absorption can be written as

\[
\mathcal{F}_{\alpha+\beta} = \left( e^{-\tau_{\text{total}}} \right) \int_{\Delta_\alpha} \int_{\Delta_\beta} e^{-\tau_{\alpha}(z_\alpha, \Delta_\alpha) + \tau_{\beta}(z_\beta, \Delta_\beta)} p(z_\alpha, \Delta_\alpha) p(z_\beta, \Delta_\beta) \, d\Delta_\alpha \, d\Delta_\beta = \int_{\Delta_\alpha} e^{-\tau_{\alpha}(z_\alpha, \Delta_\alpha)} p(z_\alpha, \Delta_\alpha) \, d\Delta_\alpha \int_{\Delta_\beta} e^{-\tau_{\beta}(z_\beta, \Delta_\beta)} p(z_\beta, \Delta_\beta) \, d\Delta_\beta = \mathcal{F}_\alpha(z_\alpha) \mathcal{F}_\beta(z_\beta) ,
\]

where \( \mathcal{F}_\alpha(z_\alpha) \) is the transmitted flux of the foreground Ly\( \alpha \) absorption at \( z_\alpha \) defined in equation (13), and \( \mathcal{F}_\beta(z_\beta) \) is the transmitted flux due solely to the Ly\( \beta \) absorption.

In Paper II we found \( \mathcal{F}_\beta = -0.002 \pm 0.020 \) in the Ly\( \beta \) GP trough of SDSS 1030+0524 at \( z_{\text{abs}} = 6 \), after correcting for the foreground Ly\( \alpha \) absorption (\( \mathcal{F}_\alpha = 0.12 \)) at \( z_{\text{abs}} = 6.0 \). Using equations (15) and (16), we find \( \Gamma_{-12} < 0.025 \) from the Ly\( \beta \) trough, an upper limit a factor of \( \sim 3 \) lower than that estimated from the Ly\( \alpha \) trough.
The oscillator strength of the Ly$\gamma$ transition is a factor of 14.3 smaller than that of Ly$\alpha$ and a factor of 2.7 lower than that of Ly$\beta$; therefore, in principle, Ly$\gamma$ can be used to put an even stronger constraint on the equivalent Ly$\alpha$ optical depth. In practice, the modeling of Ly$\gamma$ absorption is complicated due to three reasons:

1. Foreground Ly$\alpha$ absorption at $z_0 \sim 4.6$. We assume $\tau_2 \sim 0.20$, following the values in Figure 1.
2. Foreground Ly$\beta$ absorption at $z_\beta \sim 5.7$. From Figure 2, we find $\Gamma_{-12}(z = 5.7) \sim 0.12$, and the Ly$\beta$ transmission $\tau_\beta = 0.33$, using equations (13) and (15).
3. Contamination from Ly$\delta$ absorption.

In the redshift range of the GP trough $5.95 < z_{\text{abs}} < 6.16$, Ly$\gamma$ absorption covers a wavelength range $6759 \AA < \lambda < 6963 \AA$. The Ly$\delta$ absorption covers a wavelength range $6600 \AA < \lambda < 6800 \AA$, partially overlapping with the Ly$\gamma$ trough, and the Ly$\delta$ emission line is at $\lambda = 6914 \AA$. From the Keck/ESI spectrum presented in Paper II, we find that $\tau = 0.0043 \pm 0.0033$ in the wavelength range $6759 \AA < \lambda < 6963 \AA$, and $\tau = 0.0028 \pm 0.0041$ in the more restricted wavelength range $6800 \AA < \lambda < 6963 \AA$. The latter wavelength range is redward of the Ly$\delta$ GP trough. Although part of it is still bluer than the Ly$\delta$ emission, the amount of Ly$\delta$ absorption is strongly suppressed because of the proximity effect from the quasar. In the wavelength range of $6914 \AA < \lambda < 6963 \AA$, totally redward of Ly$\delta$ emission, we find $\tau = -0.0022 \pm 0.0087$.

From a combination of these transmitted flux ratio measurements, we adopt the limit $\tau_{\alpha + \beta + \gamma} < 0.007$ at $1 \sigma$ in the Ly$\gamma$ GP trough region. Correcting for the foreground Ly$\alpha$ and Ly$\beta$ absorptions, we find the transmission solely due to Ly$\gamma$ to be $\tau_\gamma < 0.10$. This transmitted flux ratio corresponds to a photoionization rate of $\Gamma_{-12} < 0.020$, comparable to the limit derived from the Ly$\beta$ GP trough and representing a factor of greater than 20 decline from the ionizing background at $z \sim 3$. The rapid decrease of $\Gamma$ at $z > 5.7$ indicates a possible sharp transition in the ionization state of the IGM (§ 5).

4. SIMULATING THE LY$\alpha$ FOREST

In this section we use an N-body simulation to simulate the absorption spectra of high-redshift quasars in the Ly$\alpha$ region, taking into account the resolution and noise properties of the observations. We fix the photoionization rate in the simulation to reproduce the observed mean transmitted flux at each redshift and compare the statistical properties of the simulated spectra, including the probability distribution function (PDF) and the threshold crossing statistics (Miralda-Escudé et al. 1996) with those of the observed spectra.

We carry out the simulations in a LCDM model, with $\Omega_m = 0.3$, $\Omega_\Lambda = 0.7$, $h = 0.65$, $\Omega_b h^2 = 0.02$, and $\sigma_8 = 0.9$. Here, $\sigma_8$ is the rms density fluctuation in 8 $h^{-1}$ Mpc spheres, chosen to reproduce the observed abundance of clusters at $z = 0$ (White, Efstathiou, & Frenk 1993; Eke, Cole, & Frenk 1996). We evolve the density fluctuations using a particle-mesh (PM) N-body code that is described in detail in Hennawi et al. (2002). This code uses a staggered mesh to compute forces on particles (Melott 1986; Park 1990), and uses the leapfrog scheme described in Quinn et al. (1997) to integrate the equations of motion. The periodic simulation cube is $L = 25$ $h^{-1}$ Mpc on a side and uses $N_p = 256^3$ particles and an $N_m = 512^3$ force mesh. We set up initial density fluctuations using the power-spectrum parameterization of Efstathiou, Bond, & White (1992) with the shape parameter $\Gamma = 0.2$ and assign initial displacements and velocities to the particles using the Zeldovich approximation. We evolve from redshift $z = 49 \rightarrow 5.45$, taking 55 equal steps in the expansion scale factor.

We use the TIPSY package to create artificial Ly$\alpha$ absorption spectra along 400 random lines of sight through the simulation cube. TIPSY calculates the local mass density ($\rho_m$) at the location of each dark matter particle using a cubic spline smoothing kernel (Hernquist & Katz 1989) enclosing 32 neighbors, and assigns it a temperature $T = T_0 \Delta^0.6$. We assumed $T_0 = 10^4$ K while creating the spectra, although the values of $\Gamma_{-12}$ we quote in this paper have been rescaled to $T_0 = 2 \times 10^4$ K. Spectra along various lines of sight are then created from this particle distribution using the algorithm described by Hernquist et al. (1996). We smooth the simulated spectra with a Gaussian filter of smoothing length $\sigma_r = 28$ km s$^{-1}$ (corresponding to a full width at half maximum of 66 km s$^{-1}$) and bin them in pixels of width 35 km s$^{-1}$, approximately twice the resolution of the Keck spectra. We then add noise to the simulated Ly$\alpha$ forest spectra. The noise in the Keck spectra is dominated by the background sky and is a strong function of wavelength around 8500 Å. In order to correctly simulate this structure in the noise, we replace the fluxes ($F$) in each simulated spectrum by the values $F = F + \sigma_{\text{obs}} G(1)$, where $\sigma_{\text{obs}}$ is the noise array in contiguous pixels starting from a random location in the Keck spectrum binned to about 35 km s$^{-1}$ and $G(1)$ is a Gaussian random deviate with zero mean and unit variance. We scale all the optical depths at each redshift so that the mean transmitted flux ratio computed from all the 400 simulated spectra match the observed values.

We compute the Ly$\alpha$ forest spectra at $z_{\text{abs}} = 5.5$, 5.7, and 6.0, and compare them with the observed Keck spectra of SDSS 1044−0130, SDSS 1306+0356, and SDSS 1030+0524 at the same wavelength ranges. Figure 3 shows the Ly$\alpha$ forest at $5.4 < z_{\text{abs}} < 5.6$ and $5.9 < z_{\text{abs}} < 6.1$ from the Keck spectra and the artificial Ly$\alpha$ absorption spectra along eight random lines of sight through the simulation. Since the simulation cube is only 25 $h^{-1}$ Mpc on a side, each simulated spectrum is of length 3506 km s$^{-1}$, 3558 km s$^{-1}$, and 3635 km s$^{-1}$ at $z_{\text{abs}} = 5.5, 5.7, \text{and } 6$, respectively. The simulated spectra are qualitatively similar to the Ly$\alpha$ forest region in the Keck spectra. Although they have the same mean transmitted flux by construction, the fluctuations in the simulated spectra arise from fluctuations in the mass density field in a LCDM model (except at $z_{\text{abs}} = 6$, where both the observed and the simulated spectra are consistent with noise). The fact that the fluctuations in the simulated spectra are qualitatively similar to those in the Keck spectra is an indication that our physical description of the Ly$\alpha$ forest as arising from photoionized gas in a low-density IGM in a LCDM model is reasonable.

Figure 4 shows the PDF of the transmitted flux ratio (Jenkins & Ostriker 1991; Miralda-Escudé et al. 1996; Rauch et al. 1997), computed using the 400 artificial spectra. The solid points in the three panels show the flux PDF measured from the Keck spectra of SDSS 1044−0125,
SDSS 1306+0356 and SDSS 1030+0524 in a velocity range of 8500 km s\(^{-1}\) in three redshift bins. We show this quantity as a function of the fraction of pixels in the spectrum that are less than the threshold flux (Miralda-Escudé et al. 1996; Weinberg et al. 1998). This statistic is analogous to the “genus curve” used to characterize the topology of the three-dimensional galaxy distribution and has the advantage that it is insensitive to the exact relation between the relative distribution of dark matter and baryons (the “biasing” of the Ly\(\alpha\) forest), as long as this relation is monotonic. We compute the cosmic variance in this statistic in a similar manner to the flux PDF. The simulated spectra reproduce this statistic of the observations very well at all three redshifts. Therefore, this modeling of the Ly\(\alpha\) forest in a LCDM model can reproduce the topology of the dark matter density fluctuations that are probed by the observed Ly\(\alpha\) forest in these quasar spectra.

5. Constraining the Epoch of Reionization

In this section, we expand the model described in § 3.1 to calculate from observations the evolution of three quantities that reflect the ionization state of the IGM: (1) the neutral fraction of the IGM, taking into account the inhomogeneity of the IGM, (2) the critical density—the density lower than which the gas is completely ionized, and (3) the mean free path of ionizing photons in the IGM. Using these calculations, we show that, as reionization progresses toward lower redshifts, the neutral fraction decreases, regions of progressively higher densities become ionized, and the mean free path of ionizing photons increases. We compare these three quantities with the results from a simulation of the cosmological reionization process by Gnedin (2000) to constrain the likely epoch of reionization (see also Gnedin 2001).

We can use equation (6) to calculate the local neutral hydrogen fraction \(f_{\text{HI}}(\Delta)\) in a region of density \(\Delta\),

\[
f_{\text{HI}}(\Delta) = \frac{n_{\text{HI}}(\Delta)}{n_{\text{HI}}(\Delta) + n_{\text{H}}(\Delta)} = \frac{n_{\text{HI}}(\Delta)}{\langle n_{\text{HI}} \rangle \Delta},
\]

where \(\langle n_{\text{HI}} \rangle\) is the average \(n_{\text{HI}}\) over the universe. We define

\[
A = 1.16 \langle n_{\text{HI}} \rangle \frac{\alpha(T)}{\Gamma},
\]

where the coefficient 1.16 comes from the helium contribution to the electron density, assuming a helium fraction \(Y = 0.24\) (Weinberg et al. 1997). Then, it can be shown that

\[
f_{\text{HI}}(\Delta) = \frac{-1 + \sqrt{1 + 4\Delta}}{1 + \sqrt{1 + 4\Delta}}.
\]

Note that for a highly ionized medium, \(f_{\text{HI}} \sim A \Delta\). We calculate the redshift evolution of the H I fraction using the density distribution of equation (13) and the results on \(\Gamma\) in Figure 2. Note that \(f_{\text{HI}}\) is directly related to \(A \sim \alpha(T)/\Gamma\). Therefore, our estimate of \(f_{\text{HI}}\) is insensitive to our assumed value for the IGM temperature \(T\), since our estimate of \(\Gamma\) is itself proportional to the assumed \(\alpha(T)\). Figure 6 shows both the volume-averaged neutral hydrogen fraction \(f_{\text{HI}}^V\) (dashed line) and the mass-averaged quantity \(f_{\text{HI}}^M\) (solid line). The volume-averaged H I fraction increases from \(1.7 \times 10^{-5}\) at \(z \sim 3.9\) to \(1.3 \times 10^{-4}\) at \(z = 5.8\)—an increase by a factor of \(\sim 8\). In the GP trough region at \(z \sim 6.05\), the upper limit on the ionization rate \(\Gamma\) from the Ly\(\alpha\), Ly\(\beta\), and...
Ly\textgamma\, transitions yields a lower limit on the neutral fraction, \(f_{\text{HI}}^{\text{V}} \gtrsim 10^{-3}\), and \(f_{\text{HI}}^{\text{M}} \gtrsim 10^{-2}\), an increase by almost two orders of magnitude from \(z \approx 4\). The mass-averaged value is larger because most \(\text{H} \text{I}\) is concentrated in dense clumps in the IGM which occupy a small volume.

The neutral hydrogen fraction depends on the detailed density distribution of the IGM. As pointed out in §3.1, \(\tau_{\text{eff}} < \langle \tau \rangle\). Therefore, simply using equation (2), which assumes a homogeneous IGM, will severely underestimate the neutral fraction. For example, at \(z = 5.8\), \(\tau_{\text{eff}} = 3.7\) (Fig. 1), and the neutral fraction derived using equation (2) is \(5.6 \times 10^{-5}\); a factor of \(\approx 2.3\) and \(\approx 70\) lower than the volume and mass-averaged neutral fractions, respectively, from equation (19). In a clumpy IGM the high-density regions quickly become opaque to all Ly\textgamma\, photons, while almost all the transmitted flux comes through the deepest voids in the IGM. We illustrate this point in Figure 7, which shows the cumulative distributions of transmitted flux as a function of density at \(z = 5.3\). At this redshift, while \(\tau_{\text{eff}} \approx 2.5\), the region with the minimum absorption in the spectrum of SDSS 1044\,0125 (Fan et al. 2000) has \(\tau < 0.5\). It represents the deepest void in the IGM, with \(\Delta \approx 0.15\) (from eq. [11]) and has a neutral fraction of \(\approx 1.2 \times 10^{-5}\), a factor of \(7\) lower than the volume-averaged neutral fraction at this redshift. It is a factor of \(\approx 50\) lower than the value one would get by directly applying equation (2) and assuming a uniform IGM, as \(\tau\) scales as \(\Delta^2\) (eq. [11]).

In all the calculations above, we have assumed that the ionizing background (and hence \(\Gamma\)) is uniformly distributed in the IGM. Figure 7 shows that the \(\Gamma\) derived from the observed transmitted flux ratio is characteristic of the background level in low density regions of the IGM, since the high-density regions contribute little to the observed flux. The observed transmitted flux (eq. [11]) is essentially a volume-averaged measurement. In a similar sense, our calculation of the volume-averaged neutral fraction should be regarded as more reliable than the mass average, as the former quantity is less affected by the local ionization state in the high-density regions.

Miralda-Escudé et al. (2000) present a semianalytic model describing the evolution of reionization in an inhomogeneous universe; reionization starts in voids and then gradually penetrates deeper into overdense regions. This process is defined by the redshift evolution of the critical density parameter \(\Delta_i\). At any redshift, regions with \(\Delta > \Delta_i\) remain neutral, while regions with \(\Delta < \Delta_i\) are mostly ionized. Here we follow Gnedin (2000) by defining \(\Delta_i\) as the critical density above which 95% of the neutral gas lies. Figure 8 shows the evolution of \(\Delta_i\) based on the neutral fraction calculated above. The quantity \(\Delta_i \approx 40\) at \(z \approx 4\), and it quickly
decreases to $D_i < 4$ in the GP trough region at $z \sim 6$. Using equation (8) of Miralda-Escudé et al. (2000), we calculate the mean free path of the ionizing photons, $\lambda_i = \lambda_0 [(1 - F_i(D_i))]^{-2/3}$, where $F_i(D_i)$ is the fraction of volume with $D < D_i$, and $\lambda_0 H = 60 \text{ km s}^{-1}$ (Miralda-Escudé et al. 2000). Figure 9 shows the evolution of the comoving mean free path with redshift. It decreases from $\sim 80 h^{-1} \text{ Mpc}$ at $z \sim 4$ to smaller than $\sim 8 h^{-1} \text{ Mpc}$ in the GP trough region at $z \sim 6$.

Figures 6, 8, and 9 clearly demonstrate that, at $z \sim 6$, the universe is much more neutral than at lower redshift ($z \sim 3$–5). At $z \sim 6$, the presence of the GP trough shows that more than 1% of the protons are in the form of $H_i$, the moderately overdense regions ($\Delta \gtrsim 4$) are still mostly neutral, and the mean free path of ionizing photons is several megaparsecs in comoving distance—a distance comparable to the correlation length of the galaxy distribution at $z = 0$. The last remaining $H_i$ in the IGM is being ionized at this epoch. In this sense, the first appearance of the complete GP trough at $z \sim 6$, marks the end of the reionization epoch.

Cosmological reionization is a complicated process. Cosmological simulations that include gas dynamics, star formation processes, atomic and molecular physics, and radiative transfer have recently begun to shed light on when and how this process happened in popular cosmological models. Gnedin (2000) divides the reionization process into three phases: the preoverlap stage in which individual $H_\text{II}$ regions begin to grow, the overlap stage in which these $H_\text{II}$ regions merge and the ionizing background in the IGM rises by a large factor, and the postoverlap stage in which the remaining $H_i$ in the high-density regions is gradually ionized. The results from the simulation of Gnedin (2000) are shown in Figures 6, 8, and 9. In the simulation, the overlap stage was assumed to be at $z \sim 7$. At this redshift, the universe goes through a near phase transition: the ionizing background and the mean free path of photons increases and the neutral hydrogen fraction decreases dramatically over a narrow redshift range. The reionization redshift can be defined in a number of ways: e.g., the redshift at which the average neutral hydrogen fraction is $\sim 50\%$, the redshift at which the $H_\text{II}$ filling factor is of order unity, the redshift at which an average region of the IGM is ionized ($\Delta_i \sim 1$), etc. However, as shown in Gnedin (2000), Chiu & Ostriker (2000), and Razoumov et al. (2001), the overlapping stage occurs over a small range in redshift, with the neutral fraction changing from 90% to less than 1% in $\Delta z < 2$ (corresponding to a time interval of 200 million yr at $z \sim 7$), while $\Gamma_{-12}$ increases from $\sim 10^{-3}$ to $\sim 10^{-1}$.

![Fig. 5.—Threshold-crossing statistic as a function of the fraction of pixels in the spectrum less than the flux threshold in the Ly$\alpha$ forest region in the Keck spectra (points) and simulated spectra (lines) at (top left) $z_{\text{abs}} = 5.5$, (top right) $z_{\text{abs}} = 5.7$, and (bottom) $z_{\text{abs}} = 6.0$. The solid line shows this statistic computed using 400 artificial spectra, and the dashed lines show the expected 1σ uncertainty in this quantity from a single spectrum of length equal to that of the real spectrum (i.e., the cosmic variance).](image-url)
Comparing the predictions from the simulation with the observations in Figures 2, 6, 8, and 9, it is evident that the redshift range $z < 5.7$ is consistent with being in the post-overlap stage of reionization, where $D_i$ increases as the remaining overdense regions in the IGM are being ionized.

The emergence of the GP trough at $z \sim 6$ and the stringent lower limits on the neutral hydrogen fraction from the presence of the Ly$\beta$ and Ly$\gamma$ GP troughs suggest the onset of a rapid transition in the ionization state of the IGM. The upper limits on the ionizing background and mean free path and the lower limits on the neutral fraction and critical overdensity we derive from the observations are all consistent with the typical values near the end of the overlap stage of reionization in the simulations. In this sense, the epoch of $z \sim 6$ is at the transition from the overlap stage to the post-
verlap stage of reionization. Even though the current observations certainly cannot probe deeper into the reionization epoch, the near phase-transition behavior of the ionization state of the IGM at \( z \sim 6 \), and the narrow redshift range over which this process occurs in cosmological reionization simulations both imply that the reionization redshift cannot be at a redshift much higher than 6.

We caution, however, that these results are based on the GP trough in only one quasar, SDSS 1030+0524. The IGM is unlikely to be reionized in an uniform manner. We have shown that at \( z \sim 6 \), the mean free path of the ionizing photons is likely to be shorter than 10 comoving megaparsecs. In this case the ionizing background could show substantial fluctuations. These fluctuations could be due to the rarity and the possible strong clustering of ionizing sources, as well as the radiative transfer effects such as shadowing of the sources. For example, in the simulation of Gnedin (2000) the ionizing background shows fluctuations of a factor of a few at \( z \sim 6 \) (his Fig. 5) in the underdense regions of the IGM (which occupy most of the volume). In order to calculate the amplitude of these fluctuations, we would need detailed simulations of the formation of the first generation of ionizing objects. This is beyond the scope of the present paper. The first generation of objects is likely to be highly biased and clustered, which would give rise to different reionization epochs along different lines of sight and in different regions of the IGM. If the universe was ionized by luminous quasars (unlikely based on current statistics on quasar evolution, Paper I), the large H\( \text{\textsc{i}} \) regions produced by them near the epoch of overlap could even lead to gaps in transmitted flux (Haiman & Loeb 1999; Miralda-Escudé et al. 2000). Thus it is quite possible that the next quasar discovered at \( z > 6 \) will show somewhat different absorption properties.

6. DISCUSSION

The presence of the GP trough in the spectrum of the highest-redshift quasar provides a first opportunity to study structure formation at high redshift and the formation of the first generation of stars and galaxies. To first order, the redshift of the reionization epoch depends on two factors: (1) the amount of small-scale power in the power spectrum of mass-density fluctuations, which determines how many halos can collapse to make stars at a certain redshift, which also determines the clumpiness of the IGM; and (2) how efficiently the collapsed halos can produce UV photons, which in turn depends on the stellar initial mass function, the efficiency of star formation, the escape factor of UV photons from protogalaxies, etc. The reionization redshift and the observed ionizing background can be used to put constraints on these factors.

In Chiu & Ostriker’s (2000) semianalytic calculations, the reionization redshift for low-density \( \Lambda \)-dominated CDM models is typically in the range of \( 7-11 \), not very far from \( z \sim 6-8 \) suggested in this paper. On the other hand, most of the non-Gaussian texture and isocurvature models have too much small-scale power and a much earlier reionization redshift \( z \gtrsim 10 \), making them difficult to reconcile with the observations presented in this paper. Barkana, Haiman, & Ostriker (2001) used the constraint on the reionization epoch to put limits on models of warm dark matter. Assuming that the reionization redshift is smaller than 7, we find that the mass of the warm dark matter particle must be smaller than 3 keV in the standard model they considered.

Paper I estimates that the SDSS will find about 20 luminous quasars in the redshift range \( 6 < z < 6.6 \) over the 10,000 deg\(^2\) survey area. High-resolution, high signal-to-noise ratio (S/N) spectra of these luminous quasars in the Lyman series absorption regions will provide valuable probes of the end of the reionization epoch and, in particular, measure the spatial inhomogeneity of the reionization process.

In order to probe deep into and beyond the reionization epoch, sources at higher redshift are needed. The Ly\( \alpha \) emission line moves out of the optical window and into the infrared at \( z \sim 6.6 \). Thus the objects become very faint in the optical wavelengths owing to Lyman series absorption by neutral hydrogen gas in the foreground at lower redshifts. Hence this is the limiting redshift for discovering sources from ground-based optical surveys. Deep and wide-field IR surveys, such as the PRIME\(^8\) mission, are required to find objects at even higher redshifts.

The stringent limits on the neutral hydrogen fraction at low redshifts from the GP test arise from the large cross-section of neutral hydrogen to Ly\( \alpha \) photons. For the same reason, the GP test quickly becomes insensitive to larger neutral hydrogen fractions. With a half hour exposure on the Keck telescope, we are able to place a lower limit on the neutral fraction \( f_{\text{\textsc{hi}}} \gtrsim 10^{-2} \). As the neutral fraction hydrogen scales with \( \tau \), or logarithmically with the flux, it is impossible to obtain a limit more than a few times stronger than the current limit for any reasonable exposure time. Therefore, the original version of the GP test cannot be used to probe deeper into the reionization epoch, for which \( f_{\text{\textsc{hi}}} \sim 1 \).

For an object observed prior to the reionization epoch, the damping wing of the GP trough arising from the large GP optical depth of the neutral medium will extend into the red side of the Ly\( \alpha \) emission line (Miralda-Escudé 1998). The depth and extent of this damping wing can, in principle, be used to measure the neutral fraction and the redshift of the reionization epoch. However, this GP damping wing test cannot be applied to luminous quasars because of the proximity effect from the quasar itself. As shown by Madau & Rees (2000) and Cen & Haiman (2000), the quasar will ionize a H\( \text{\textsc{i}} \) region around it if the bulk of the IGM is still neutral, resulting in a line profile very similar to the case where the IGM has already been ionized. An alternative is to search for lower luminosity quasars (Stern et al. 2000) and galaxies (Hu et al. 1999; Rhoads & Malhotra 2001) at high redshifts by either applying the color dropout technique to deep multiband imaging data or by searching in regions around foreground galaxy clusters where background objects might be amplified by gravitational lensing (e.g., Ellis et al. 2001). The highly magnified sources may be bright enough to allow high-S/N spectroscopy, necessary for accurately measuring the profile of the damping wing.

Another promising opportunity for probing the reionization epoch arises from the observations of IGM absorption in the afterglow of high-redshift gamma-ray bursts (GRBs; Loeb 2001). GRBs are transient phenomena with no immediate impact on the surrounding IGM. The time dilation at high redshift leads to a longer duration for the afterglows.

\(^8\) See http://prime.pha.jhu.edu/index.html.
making the identification and follow-up observations easier. Moreover, if high-redshift GRBs are associated with the collapse of massive stars, their evolution is likely to be similar to those of star-forming galaxies, which show slower decline at high redshift (Steidel et al. 1998) than the rapid decline in the number density of quasars (e.g., Fan et al. 2001b).

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