A "MINIHALO" MODEL FOR THE LYMAN LIMIT ABSORPTION SYSTEMS AT HIGH REDSHIFT

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Received 1997 August 29; accepted 1997 December 8; published 1998 February 2

ABSTRACT

We propose that a large fraction of QSO Lyman limit absorption systems (LLSs) observed at high redshift \((z \gtrsim 3)\) originate from gas trapped in small objects, such as minihalos, that form prior to reionization. In the absence of a strong UV flux, the gas is predominantly neutral and may form clouds with H\textsc{i} column density \(N_{\text{H}\textsc{i}} \gtrsim 10^{18} \text{ cm}^{-2}\). Owing to their high densities and high H\textsc{i} column densities, these clouds are not destroyed by the onset of the UV background at a later time. Thus, if not disrupted by other processes, such as mergers into larger systems or "blow away" by supernovae, they will produce LLSs. We show that the observed number density of LLSs at high redshifts can be well reproduced by the survived "minihalos" in hierarchical clustering models such as the standard cold dark matter model. The number density of LLSs in such a population increases with \(z\) even beyond the redshifts accessible to current observations and dies off quickly at \(z \lesssim 2\). This population is distinct from other populations because the absorbing systems have small velocity widths and a close to primordial chemical composition. The existence of such a population requires that the reionization of the universe occurs late, at \(z \lesssim 20\).

Subject headings: cosmology: theory — galaxies: formation — quasars: absorption line

1. INTRODUCTION

Lyman limit absorption systems (LLSs), observed as breaks in quasar spectra at the rest-frame wavelength of the Lyman continuum edge, are produced by gaseous clouds optically thick to Lyman continuum photons. These systems can be observed over a wide range of redshift, from \(z \sim 0\) to approximately that of the highest redshift quasars. Thus, such systems can be used to probe the gas in the universe at various redshifts. Because of the highest redshift quasars, one might conclude that the LLSs are associated with galaxies, in which case the number density and sizes of galaxies are expected to cause the LLSs observed at high redshifts. However, in a hierarchical model of structure formation, in which the number density of LLSs at high redshift is expected to be of the order of one for each line of sight, the number density of LLSs increases with \(z\) even beyond the redshifts accessible to current observations.

It is possible that some of the LLSs at high redshifts are associated with galaxies, but at \(z \lesssim 3\) it would mean that most gas has not yet been assembled into large galaxies (Mo & Miralda-Escudé 1996). It is therefore unlikely that LLSs observed at high redshifts are caused by galactic populations.

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2. A NEW POPULATION OF LYMAN LIMIT SYSTEMS

For simplicity we assume the universe to be flat with the cosmic density parameter $\Omega = 1$. The Hubble constant is written as $H_0 = 100 \, h \, \text{km s}^{-1} \, \text{Mpc}^{-1}$.

2.1. Physical Properties

Prior to the epoch of reionization, the physical properties of cosmological structures can be estimated from the spherical collapse model (see, e.g., Padmanabhan 1993). The gas in a collapsing dark matter perturbation is shock-heated to the virial temperature

$$T_{\text{vir}} = \frac{\mu V^2}{3k} \approx 4.0 \times 10^3 \, \text{K} \left( \frac{V}{10 \, \text{km s}^{-1}} \right)^2 \left( \frac{\mu}{m_p} \right), \quad (1)$$

where $V$ is the three-dimensional virial velocity of the halo, $m_p$ is the proton mass, and $\mu$ is the mean mass per particle.

The total $H$ number density in a dark matter halo that collapses at a redshift $z$ with an overdensity $\delta$ at virialization reads

$$n_{\text{vir}} = \frac{\Omega_b \rho_{\text{crit}} \delta}{m_p} (1 + z)^3 \approx 9.7 \times 10^{-3} \times$$

$$\times \left( \frac{\Omega_b h^2}{0.0125} \right) \left( \frac{1 + z}{10} \right)^3 \left( \frac{\delta}{18 \pi^2} \right) \, \text{cm}^{-3}, \quad (2)$$

where $\rho_{\text{crit}}$ denotes the closure density. Assuming that dark halos are singular isothermal spheres with density profile $\rho(r) \propto r^{-2}$, the virial radius $r_{\text{vir}}$ defined as the radius within which the mean mass overdensity is $18\pi^2$, can be written

$$r_{\text{vir}} = \frac{V t}{2\pi} \approx 350 \, \text{pc} \, h^{-1} \left( \frac{V_c}{10 \, \text{km s}^{-1}} \right) \left( \frac{1 + z}{10} \right)^{-3/2}, \quad (3)$$

where $t = (2/3H_0)(1+z)^{-3/2}$ is the cosmic time at $z$.

For halos with $T_{\text{vir}} \approx 10^4 \, \text{K}$ ($V_c \approx 15 \, \text{km s}^{-1}$), the gaseous component cannot be collisionally ionized and will remain neutral. Owing to the high collapse redshifts ($z > 5$) considered here, the recombination timescale in minihalos is shorter than the Hubble time, and the gas in clouds with $V_c \approx 15 \, \text{km s}^{-1}$ before reionization will also be neutral shortly after their formation. As a result, the column density of neutral hydrogen atoms, $N_{\text{HI}}$, is given by

$$N_{\text{HI}} \approx 2.1 \times 10^{19} \left( \frac{V_c}{10 \, \text{km s}^{-1}} \right)^4$$

$$\times \left( \frac{1 + z}{10} \right)^{3/2} \left( \frac{\delta}{18 \pi^2} \right) \left( \frac{\Omega_b h^2}{0.0125} \right) \, \text{cm}^{-2}. \quad (4)$$

The corresponding total mass of such a halo is

$$M = \frac{V t^2}{2\pi G} \approx 7.9 \times 10^6 \left( \frac{V_c}{10 \, \text{km s}^{-1}} \right)^5$$

$$\times \left( \frac{1 + z}{10} \right)^{-3/2} h^{-1} M_\odot. \quad (5)$$

It is clear that the optical depth at the photoionization threshold,

$$\tau = \sigma(n_{\text{HI}}) N_{\text{HI}} \approx 130 \left( \frac{V}{10 \, \text{km s}^{-1}} \right)^{1/2}$$

$$\times \left( \frac{\delta}{18 \pi^2} \right) \left( \frac{\Omega_b h^2}{0.0125} \right) \, \text{cm}^{-2}, \quad (6)$$

exceeds unity for a wide range of collapse redshifts and virial velocities.

To determine what happens to the clouds once the UV background switches on, we analyze the equilibrium equation of the ionized fraction for a pure hydrogen gas, $x = n_e/n$, where $n_e$ and $n$ denote the free electron and total baryonic number density, respectively. This fraction is given by

$$0 = k_{\text{ps}} n(1-x) - k_{\text{rec}} n^2 x^2, \quad (7)$$

where $k_{\text{ps}} = 4\pi [\nu_{\text{HI}} \sigma(H)] \exp[-N_{\text{HI}} \sigma(v) \sigma(v) \, dv$ and $k_{\text{rec}}$ denote the rate coefficients for photoionization of neutral hydrogen and radiative recombination to $H$ 1, respectively. Motivated by the indirect measurements from the proximity effect (see Giallongo et al. 1996 and references therein), we assume the background UV flux to be given by a power law

$$J(\nu) = 10^{-21} J_{21} \left( \frac{h\nu}{1 \, \text{ryd}} \right)^{-\alpha} \, \text{ergs cm}^{-2} \, (\text{Hz s sr})^{-1}, \quad (8)$$

where $h$ is Planck’s constant and $h\nu_{\text{HI}} = 1 \, \text{ryd} \approx 13.6 \, \text{eV}$ is the $H_1$ ionization threshold. In the optical thick limit, the photoionization rate coefficient can be approximated by a power law $k_{\text{ps}} \approx J_{21} k_{\text{ps}}^{\alpha} |N_{\text{HI}}(1-x)\|^{\beta}$, where $N_{\text{HI}}$ is the total column density of hydrogen nuclei. The exponent $\beta$ has a non-trivial dependence on $\alpha$ and is determined numerically to vary from 1.35 to 1.69 for $1 < \alpha < 2$. Inserting this approximate expression for the photoionization rate into the equilibrium equation for the free electron fraction (eq. [7]) yields the implicit solution for $x$:

$$f(x) = \frac{x^2}{(1-x)^{\alpha-\beta}} = \frac{J_{21} k_{\text{ps}}^{\alpha} N_{\text{HI}}^{-\beta}}{nk_{\text{rec}}} = C, \quad (9)$$

which for $\beta = 0$ reduces to the more familiar optical thick case. We adopt $\alpha = 1.8$ for the spectral index, with which the photoionization rate can be approximated by $k_{\text{ps}} \approx 4.1 \times 10^{-17} J_{21} (N_{\text{HI}} / 10^{20} \, \text{cm}^{-2})^{-\beta} \, \text{s}^{-1}$ with $\beta = 1.6$.

The right-hand side of equation (9) has a maximum at $x = x_{\text{max}} \approx 0.8$ with a value of $f(x_{\text{max}}) \approx 0.25$. The existence of this maximum shows that ionization equilibrium cannot be achieved for $J_{21} k_{\text{ps}}^{\alpha} N_{\text{HI}}^{-\beta} > f(x_{\text{max}})$. In this case, photoionization overcomes recombination and ionizes the gas until a new, highly ionized, equilibrium state is reached at column densities $< 10^{17} \, \text{cm}^{-2}$. As discussed in Mückel & Kates (1997), the evolution of such systems from one equilibrium state to another can be rapid. The recombination rate suitable for an equilibrium temperature of $10^4 \, \text{K}$ is $k_{\text{rec}} \approx 2.6 \times 10^{-13} \, \text{cm}^3 \, \text{s}^{-1}$ (Ferland et al. 1992). Thus, for a given $J_{21}$, the maximum of $f(x)$ can be used to derive a critical total hydrogen column density $N_{\text{HI}}^{\text{crit}}$. Note that this column density is calculated at the edge of the cloud and is approximately equal to $r_{\text{vir}} n_{\text{HI}}$ (see eqs. [2] and [3]). We can therefore define a minimum collapse redshift, $z_{\text{crit}}$, for halos of a given virial velocity, so that they will remain optically thick if they form at $z > z_{\text{crit}}$. This min-

Note that the case B recombination rate is used for the optical thick case considered here.
Um collapse redshift can be written as

$$1 + z_{\text{crit}} \approx 7.2 \left( \frac{\Omega_\Lambda}{0.5} \right)^{5/27} \left( \frac{V}{10 \text{ km s}^{-1}} \right)^{-8/27} \times \left( \frac{\delta}{18\pi^2} \right)^{-1/27} \left( \frac{h}{0.5} \right)^{-2/3}, \quad (10)$$

As we will show below, minihalos (with $V \approx 10$ km s$^{-1}$) in current hierarchical models have typical collapse redshifts higher than this critical value; they can therefore retain their high initial H I column densities.

To check the robustness of the above derivation for the critical collapse redshift, we have integrated the time-dependent chemistry and cooling model presented in Abel et al. (1997). We found good agreement between the numerical results and the analytic expression of the critical redshift given by equation (10).

The density of gas contained in minihalos considered here is quite high because of their high collapse redshifts. The recombination timescale for the gas is short, about 0.1 Myr, which makes them stable against photoionization from internal stellar sources. On the other hand, only a few supernova explosions may be able to blow out most of the gas from the halos (see Ciardi & Ferrara 1997 and references therein) if star formation can happen in them. However, systems that form in halos with $V \lesssim 15$ km s$^{-1}$ at $z \approx 15$ are not able to cool since H$_2$ formation is inefficient in these halos at such redshifts (see Tegmark et al. 1997 and Abel et al. 1998). Also, the UV flux of the first structures in the universe may lead to H$_2$ photodissociation prior to H I reionization (Haiman, Rees, & Loeb 1997), inhibiting line cooling by H$_2$. Furthermore, H I line cooling is not efficient because of the low virial temperature ($\lesssim 10^4$ K). Thus, these systems may not form stars and hence are not subject to “blow away.” The gas can only be gravitationally confined in the dark matter halo if the typical particle velocity is less than the escape velocity of the halo. The mean particle velocity of hydrogen atoms in a Maxwellian distribution is $V_n(T) = (3kT/m_p)^{1/2}$ for a neutral gas, and the escape velocity is given roughly by $V_{\text{esc}} \approx \sqrt{3}V_n$. Hence, only systems with $V \approx V_n(10^4$ K)/$\sqrt{3} \approx 9$ km s$^{-1}$ are stable against evaporation. These considerations suggest that minihalos with $V = 10$–20 km s$^{-1}$ are most likely to produce LLSs.

### 2.2. The Predicted Number Density of LLSs

The typical collapse redshift of halos with $V \approx 30$ km s$^{-1}$ in current cosmogonies such as the standard CDM model exceeds the critical redshift given in equation (10). Such small halos are also abundant at high redshift in these models. It is therefore possible that minihalos contribute a substantial part of the total number of LLSs at high redshifts. Here we examine this possibility.

We assume that minihalos that form at $z > z_{\text{crit}}$ and do not merge into larger systems by a redshift $z_*$ will contribute to the absorption cross section at $z_*$. The conditional probability, $p(M_2, z_2 | M_1, z_1) dM_2$, that a halo of mass $M_1$ selected at redshift $z_1$ will have merged to form a halo of mass between $M_1$ and $M_1 + dM_2$ at a later redshift $z_2$ can be calculated from the extended Press-Schechter formalism (Bond et al. 1991; Bower 1991). As in Lacey & Cole (1994), halos are assumed to be destroyed by redshift $z_*$ if they have merged into halos of twice their initial mass by that redshift. The fraction of minihalos that form at $z_1$ and survive merging until $z_*$, $f_s$, can then be obtained from

$$f_s(M_1, z_1, z_*) = 1 - \int_{2M_1}^{\infty} p(M_2, z_2 | M_1, z_1) dM_2. \quad (11)$$

The absorption cross section of a minihalo is assumed to be spherical with radius given by equation (3). The result for a standard CDM model normalized to $\delta = 0.7$ is given for three different bins of virial velocities in Figure 1. In the calculation $z_1$ is chosen as the redshift at which the comoving number density of halos predicted by the Press-Schechter formalism peaks. This corresponds to the collapse redshift of a $\sim 1.5$ σ peak. In all cases, $z_* > z_{\text{crit}}$.

Clearly, the total number of systems with $10$ km s$^{-1} \lesssim V \lesssim 15$ km s$^{-1}$ is sufficient to explain the observed number of LLSs at $z \approx 3$. This is consistent with the observational evidence that LLSs at low redshift are associated with galaxies (see Steidel et al. 1996). As shown in Mo & Miralda-Escudé (1996), clouds pressure-confined in galactic halos can indeed give a sufficiently large cross section to explain the number density of LLSs at $z \approx 2$. Note that for current low-$\Omega$ cosmogonic models having smaller power on small scales than the standard CDM model, the number density of minihalos may be reduced. However, the predicted number of absorption systems per unit redshift may not be reduced substantially, because it depends also on $dl/dz$ (the proper distance per unit redshift) and the virial radius; both are larger in a low-$\Omega$ model. The dependence on cosmological parameters cannot be examined in detail until more detailed hydrodynamical calculations are able to determine the physical sizes of the absorbers.

It is important to realize the difference between our model and the minihalo model proposed by Rees (1986) for Ly$\alpha$ forest systems. In the model of Rees, minihalos are optically thin, highly ionized, and have an H I column density of $\lesssim 10^{17}$ cm$^{-2}$, characteristic for Ly$\alpha$ forest systems. This assumption is correct for minihalos that collapse after reionization from an already
ionized intergalactic medium (IGM) at redshifts $z < z_{\text{crit}}$ (see eq. [10]).

3. DISCUSSION

As shown in Thoul & Weinberg (1996), after reionization the collapse of gas can only happen in halos with $V_c \geq 30$ km s$^{-1}$, and so minihalos (with $V_c \leq 30$ km s$^{-1}$) that form after reionization are not expected to give a significant contribution to the LLSs. In contrast, prior to reionization the collapse of gas can even happen in halos with $V_c$ as small as about 0.1 km s$^{-1}$($1 + z_{\text{coll}}$). Our model for the LLSs envisages a hierarchical cosmology where a large number of minihalos (with $V_c \leq 30$ km s$^{-1}$) can form prior to the epoch of reionization. We show that the gas trapped in such minihalos is, because of its high density, able to withstand photoionization by the UV background set on at a later time and reaches thermal and chemical equilibrium at a temperature $\sim 10^4$ K. By including these systems, the discrepancy in the number of LLSs between the observational result and earlier theoretical predictions based on CDM models can be reconciled. The model predicts the number density of LLSs to keep increasing toward higher redshifts.

The minihalo population is assumed to form at a time when the gas in the IGM is largely primordial. As a result, the absorbing gas in these systems should also have a close to primordial composition. Although later accretion of gas from the enriched IGM may increase the metallicity, we expect the metallicity of such LLSs to be low.

Minihalos with $10$ km s$^{-1} \leq V_c \leq 15$ km s$^{-1}$ are expected to dominate the total cross section of LLSs at high redshifts, because (1) they can form in the low-pressure IGM prior to reionization, (2) they are stable against evaporation, (3) they cannot cool by H$_2$ or H$_i$ to form stars and so are not subject to supernova explosion or internal photoionization, and (4) their total cross section is large in hierarchical models. Thus, the minihalo population of the LLSs should be observed to have a narrow range of velocity widths (with H$_i$ Doppler parameters of $\leq 20$ km s$^{-1}$). This can be tested by high-resolution spectroscopy of LLSs at high redshifts. The two LLSs at redshifts 3.32 and 2.80 studied by Tytler, Burles, & Kirkman (1996) using the HIRES spectrograph on the KECK telescope have Doppler parameters and velocity offsets between individual components less than 20 km s$^{-1}$. This is consistent with our model. Obviously, more high-resolution data at high redshifts are needed to constrain the fraction of LLSs in the minihalo population. The sizes for the minihalo absorbers are typically about 1 kpc, which may be tested by the spectra of gravitationally lensed quasar pairs at high redshifts.

Much theoretical work remains to be done to model the hydrodynamic and radiative transfer processes so as to obtain more accurate predictions for the properties of the minihalo Lyman limit systems proposed here.

We thank Jordi Miralda-Escudé, Avi Loeb, Karsten Jedamzik, Sandra Savaglio, Paul Shapiro, Andrea Ferrara, and Simon White for comments on an earlier draft of this Letter. We are especially grateful to Martin Haehnelt for many discussions and his constructive criticism on this work. T. A. acknowledges support from NASA grant NAG5-3923.

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