HEATING AND COOLING OF THE EARLY INTERGALACTIC MEDIUM BY RESONANCE PHOTONS

LEONID CHUZHoy AND PAUL R. SHAPIRO

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ABSTRACT

During the epoch of reionization a large number of photons were produced with frequencies below the hydrogen Lyman limit. After redshifting into the closest resonance, these photons underwent multiple scatterings with atoms. We examine the effect of these scatterings on the temperature of the neutral intergalactic medium (IGM). Continuum photons, emitted between the Lyα and Lyγ frequencies, heat the gas after being redshifted into the H Lyα or D Lyβ resonance. By contrast, photons emitted between the Lyγ and Lyman limit frequencies produce effective cooling of the gas. Prior to reionization, the equilibrium temperature of ~100 K for hydrogen and helium atoms is set by these two competing processes. At the same time, Lyβ resonance photons thermally decouple deuterium from other species, raising its temperature as high as 10^4 K. Our results have important consequences for the cosmic 21 cm background and the entropy floor of the early IGM, which can affect star formation and reionization.

Subject headings: diffuse radiation — intergalactic medium — radiation mechanisms: general

1. INTRODUCTION

During the epoch of reionization, UV photons between the Lyα and Lyman limit frequencies are produced in abundance. Unlike the ionizing photons above the Lyman limit, whose mean free path in the neutral IGM is quite short, photons between Lyα and the Lyman limit scatter only after redshifting into one of the atomic resonances, which allows them to move freely through cosmological distances. Eventually, these scattering photons exchange energy with atoms, thereby providing a mechanism for raising or lowering the gas temperature far from any radiation source.

Several authors previously considered this process. Madau et al. (1997) estimated that photons lose about 1% of their energy to the gas as they redshift through Lyα resonance. They found accordingly that, if Lyα scattering controls the H i hyperfine level population, as required to produce a 21 cm background distinct from the cosmic microwave background (CMB), recoil heating sufficed to raise T_IGM above T_{CMB} and to move the 21 cm background into emission. Later, however, Chen & Miralda-Escude (2004) showed that the inclusion of atomic thermal motion (neglected by Madau et al. 1997) in the analysis reduces the heating rate by at least 3 orders of magnitude. Furthermore, they showed, while photons emitted between Lyα and Lyβ frequencies (called by them the “continuum photons”) heat the gas when redshifted into Lyα resonance, those injected directly into Lyα by the cascade from higher resonances (called the “injected photons”) cool the gas. With increasing gas temperature, the efficiency of “continuum photons” as heaters declines, while that of “injected photons” as coolers rises. Therefore, if, as these authors claimed, the numbers of these two types of photons were similar, then cooling would prevail at temperatures above 10 K. If so, the 21 cm background produced by Lyα pumping would have been in absorption, unless some other mechanism, such as photoelectric heating by a strong X-ray background, can raise the IGM temperature above T_{CMB} without fully ionizing it.

However, two important physical mechanisms were missing in previous calculations. The first is the forbidden transition from the 2s to 1s level in hydrogen. The cascade that follows the absorption of most photons that redshift into high atomic resonances (Lyβ, Lyγ, Lyδ, etc.), proceeds preferentially not via the 2p level (whose decay to 1s produces a Lyα “injected photon”), but via the 2s level instead. This significantly reduces the number of injected photons, raising the equilibrium temperature to about 100 K. The second is the impact of deuterium. Despite its very low abundance, deuterium makes an important contribution to the gas heating by its interaction with Lyβ photons. At high radiation intensities, the D Lyβ resonance becomes more important for heating the gas than H Lyα. In addition, scatterings by D atoms reduce the number of injected photons for hydrogen, thus decreasing the cooling rate.

In § 2 and § 3 we calculate the effect of resonant scattering on H and D atoms, respectively. In § 4 and § 5 we discuss the impact of this scattering on the thermal evolution, 21 cm signal and reionization history of the IGM.

2. HYDROGEN RESONANCES

When photons redshift through the Lyα resonance, multiple scatterings affect their spectrum. The intensity J(ν) varies with frequency in the neighborhood of the resonance at ν_ν_α according to the following (Chuzhoy & Shapiro 2005)

\[ J(\nu) = J(0)e^{-2\pi x^2/3a-2\eta x}, \]

for injected photons and x > 0, otherwise

\[ J(x) = 2\pi J_0 a^{-1} e^{-2\pi x^2/(3a-2\eta x)} \int_{-\infty}^{x} e^{2\pi z^2/3a+2\eta z^2} dz, \]

where \( a = A_{21}(2kT_H/mc^2)^{-1/2}/4\pi \nu_\alpha \), T_H is the hydrogen kinetic temperature, A_{21} is the Einstein spontaneous emission coefficient of the Lyα transition, x = ν/ν_α - 1/(2kT_H/mc^2)^{1/2} is the frequency distance from line center divided by the thermal width of the resonance, \( \gamma^{-1}(1 + 0.4/T) \) is \( \tau_G \approx 1 \) is the Gunn-Peterson optical depth, and J_0 is the UV intensity far away from the resonance. The recoil parameter, \( \eta \), equals

\[ \frac{h\nu_\alpha}{(2kT_H/mc^2)^{1/2}} \frac{1 + 0.4/T_H}{1 + 0.4/T_H} \]

1 McDonald Observatory and Department of Astronomy, University of Texas at Austin, Austin, TX; chuzhoy@astro.as.utexas.edu; shapiro@astro.as.utexas.edu.

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where $T_s$ is the spin temperature of the 21 cm hyperfine transition, which equals $T_\text{H} (T_\text{CMB})$ at high (low) radiation intensities. The intensity at $x = 0$, $J(0)$, is given by

$$J(0) = \frac{\pi \zeta (J_{1/3}(\zeta) - J_{-1/3}(\zeta))}{\sqrt{3}} + F_2 \left(1 - \frac{1}{3} \cdot \frac{2}{3} - \frac{\zeta^2}{4}\right), \quad (3)$$

where $\zeta = (16 \pi^2 a^6 \pi\gamma) \sqrt{2}$, $F_2$ is a hypergeometric function, and $J_{1/3}$ and $J_{-1/3}$ are the Bessel functions of the first kind.

Since photons to the red (blue) of the resonance scatter preferentially off atoms moving toward (away) from them, the average energy an atom gains from a scattering photon depends on its frequency

$$\Delta E(x) = \frac{(\hbar \nu)^2}{mc^2} \left[ 1 - \frac{kT_H}{\hbar} \phi'(x) \right], \quad (4)$$

where $\phi(x)$ is the normalized scattering cross section (Chuzhoy & Shapiro 2005). The total energy gain of the gas from each photon as it passes through the resonance is determined by the radiation intensity, the probability of scattering, and the average gain at each frequency

$$\Delta E_{\text{tot}} = \int \frac{J(x)}{J_0} \Delta E(x) \phi(x) \, dx. \quad (5)$$

Figure 1 shows the heating and the cooling rates obtained by numerical integration of equation (5).\footnote{Results are virtually insensitive to the choice of Lorentz vs. Voigt profile for $\phi(x)$.} For $T_\text{H} \gtrless 100 \text{ K}$ and the range of optical depths corresponding to the mean density IGM at 10 $< z < 30$, $\Delta E_{\text{tot}}$ can be well approximated by the asymptotes

$$\Delta E_c/k \sim 0.37 T_\text{H}^{1/3} (1 + z) \text{ K} \quad (6)$$

for continuum photons and

$$\Delta E_i/k \sim -0.3 T_\text{H}^{1/3} (1 + z)^{1/2} \text{ K} \quad (7)$$

for injected photons. The total heating/cooling rate is

$$H_\alpha = N_\alpha \left( \Delta E_c + \frac{J_i}{J_c} \Delta E_i \right), \quad (8)$$

where $N_\alpha$ is the number of the photons that pass through the Ly$\alpha$ resonance per H atom per unit time.

The continuum photons (i.e., those that are redshifted into the Ly$\alpha$ resonance) can be produced in two ways. Most originate as photons emitted between Ly$\alpha$ and Ly$\beta$, which eventually redshift into the Ly$\alpha$ resonance. The rest come from higher frequency photons that are absorbed by D atoms and cascade into the D Ly$\alpha$ resonance, from which they then redshift into the H Ly$\alpha$ resonance. The injected photons (i.e., those injected directly into the Ly$\alpha$ resonance) come from the photons emitted between Ly$\gamma$ and the Lyman limit. Most of the latter photons, when redshifted into the closest resonance and absorbed by H atoms, produce a cascade to the 2s level, from which an electron decays to the ground level by emitting two photons below Ly$\alpha$ as also noted by Hirata (2006) and Pritchard & Furlanetto (2006). Another fraction are absorbed by D atoms and either destroyed or converted to D Ly$\alpha$ photons, as previously explained. The rest, after redshifting into the closest H resonance and producing a cascade to the 2p level, become the injected photons. Therefore, the ratio of injected and continuum photons, is

$$\frac{J_i}{J_c} = \sum_{n=3}^{\infty} \int_{\nu_{\ell}}^{\nu_{n+1}} J_d(\nu)p_n(1 - f_{D,n}) d\nu$$

$$\times \left[ \int_{\nu_{n+1}}^{\nu_{n+1}} J_d(\nu) d\nu + \sum_{n=3}^{\infty} \int_{\nu_{n+1}}^{\nu_{n+1}} J_d(\nu)f_{D,n} d\nu \right]^{-1}, \quad (9)$$

where $\nu_n$ is the frequency of $np \rightarrow 1s$ transition [i.e., $\nu_n = \nu_{1s}(1 + n^2)$, where $\nu_{1s} = 13.6 \text{ h}^{-1} \text{ eV}$ is the Lyman limit frequency], $p_n$ is the fraction of all cascades from level $np$ that go through $2p$ (0.26 for $n = 3$ and 4, respectively, and roughly 1/3 for $n > 4$), $J_d(\nu)$ is the spectral profile of the radiation source, and $f_{D,n}$ is the ratio of the cascades from level $np$ occurring in deuterium and hydrogen. It can be shown that, since the deuterium optical depth $\tau_{D,n}$ in the $np \rightarrow 1s$ transition was low during reionization, $f_{D,n} \approx 1 - e^{-\tau_{D,n}}$ (1.1 $\mu$eV), where $f_{\text{esc}}$ is the probability for the electron in level $np$ to decay directly to $1s$. For a hot source with surface temperature $T_{\text{surf}} \gg 10^5 \text{ K}$, $J_d(\nu) \propto \nu^2$ between Ly$\alpha$ and Lyman limit, and $J_{1s}/J_{2s} \approx 0.17$ (the exact value depends on the optical depth of deuterium, which is redshift dependent). Colder sources would produce a lower $J_{1s}/J_{2s}$ ratio, but unless the spectrum color temperature is below $\approx 5 \times 10^4 \text{ K}$, $J_{1s}/J_{2s} > 0.1$.

Combining equations (6) and (7), we find that for the mean IGM, the equilibrium temperature at which cooling by injected photons balances heating by continuum photons is

$$T_{\text{eq}} \approx 130 \text{ K} \left( \frac{1 + z}{10} \right)^{1/2} \left( \frac{J_1/J_{2s}}{0.15} \right)^{-3/2}. \quad (10)$$

This greatly exceeds the mean temperature prior to reionization, if only adiabatic cooling occurs once the IGM decouples from the CMB.

3. Deuterium resonances

Unlike hydrogen, the most important resonance for deuterium is not Ly$\alpha$ but Ly$\beta$ (Chuzhoy & Shapiro 2005). As photons redshift through the D Ly$\beta$ resonance, a significant fraction of them
are destroyed via absorption and cascade. Prior to reionization \((z \approx 10)\), the photon destruction probability is above 0.2. Therefore, the blue wing of the resonance is significantly higher than the red wing, so that the radiation color temperature around the resonance is negative. Because of their negative color temperature, the D Ly\(\beta\) photons are much more efficient than the Ly\(\alpha\) photons at heating the gas. If the velocity distribution of deuterium atoms is Maxwellian, then the radiation intensity around the Ly\(\beta\) resonance varies according to

\[
J(x) - \gamma_D J'(x) = 0.5 f_{\text{rec.}} \int_{-\infty}^{\infty} \text{Erfc}(\max(|x|, |x|)) J(x') dx',
\]

where \(\gamma_D^{-1} = T_{D,1} \approx 3 \times 10^3 (1 + z)^{3/2} (n_D/n_H)\) (Chuzhoy & Shapiro 2005). Solving equations (5) and (11) numerically, we find that for neutral IGM between \(z = 10\) and 20 the average energy each Ly\(\beta\) photon transfers to the gas is

\[
\Delta E_{\text{tot}} / k \approx 0.012 T_{D}^{1/2} \left(1 + z\right)^{3/2} \left(\frac{15}{32}\right) \text{K},
\]

where \(T_D\) is the deuterium kinetic temperature.

Physical insight into the above result can be gained from the following argument. Photons are scattered by atoms in whose rest frame they are close to resonance [i.e., \(\nu(1 - \nu/c) \approx \nu_{Ly} \rho\), where \(\nu\) is the velocity of an atom]. Thus, photons in the red (blue) wing of the resonance are scattered preferentially by atoms moving toward (away) from them. When the spectrum around the resonance is flat, then to first approximation, the energy gain of atoms from photons in the blue wing is compensated by energy loss to photons on the red wing. However, when the spectrum is tilted, as in our case, each extra photon on the blue wing adds to the gas \(\Delta E_{\text{tot}} / k \approx (\hbar / c / T_{D}) \nu_{\text{thermal}} c \approx 0.05 T_{D}^{1/2} \text{K}\), comparable to what we got in equation (12). We note that, unlike H Ly\(\alpha\) photons, D Ly\(\beta\) photons become more efficient at heating when the temperature rises.

Since radiation affects H and D atoms differently, we now have to make separate estimates of their respective temperatures. From equation (12) we find that the heating rate per D atom is

\[
H_{\beta} \approx 4 \times 10^{-28} \text{ergs s}^{-1} \left(\frac{n_D/n_H}{2} \times 10^{-15}\right)^{-1} \left(\frac{H_{\beta}}{10^{14}}\right) ^{3/2} T_{D}^{1/2},
\]

where \(H_{\beta}\) is the number of photons passing through the Ly\(\beta\) resonance per H atom per second. Elastic collisions with H atoms, and to a smaller extent with He atoms, make D atoms lose energy at a rate

\[
L_D \approx k(T_D - T_H) \left(\frac{3kT_D}{2m_D} + \frac{3kT_H}{2m_H}\right)^{1/2} n_D \sigma_{HD}
\]

\[
= 1.3 \times 10^{-31} \text{ergs s}^{-1} (T_D - T_H) \left(\frac{T_D}{2} + T_H\right)^{1/2} \left(\frac{\Omega_{H}}{0.02}\right) \left(\frac{1 + z}{10}\right)^{3} \left(\frac{\sigma_{HD}}{10^{-15} \text{cm}^2}\right)^{-1},
\]

where \(\sigma_{HD}\) is the momentum transfer cross section for collisions between H and D atoms. When the radiation intensity is high, so that \(T_D \gg T_H\), \(T_D\) is set by the balance between heating and cooling rates, \(L_D = H_{\beta}\).

FIG. 2.—Thermal evolution of H (solid lines) and D (dashed lines) atoms for \(n_D = 10^{-15} \text{cm}^2\) and \(N_{\gamma,H} (z = 11) = 200\). The convolving radiation intensity is assumed to grow as \(e^{-z^2/\Delta z^2}\), where \(\Delta z = 1, 2, 3\) correspond to lower, middle, and upper lines.

4 Some element of uncertainty is present in the above equation due to the absence of experimental measurements of \(\sigma_{HD}\). The collisional cross sections of atoms belonging to different species can be roughly determined by the size of the atom, which for D and H atoms gives \(\sigma_{HD} = 0.35 \times 10^{-15} \text{cm}^2\) (Clementi et al. 1963; Slater 1964). However, since D and H atoms energy levels are nearly degenerate it is possible that \(\sigma_{HD}\) is significantly larger.
higher. When gas collapses out of the IGM, the entropy cannot decrease (without radiative cooling), so its overdensity inside virialized halos cannot exceed \((T_{\text{vir}}/T_0)^{3/2}\), where \(T_0\) and \(T_{\text{vir}}\) are the temperatures before (after) collapse, respectively. Consequently, for minihalos with \(T_{\text{vir}} \lesssim 4000\) K that form out of gas preheated to 100 K, the post-collapse overdensity is less than \(~200\). Minihalos with such lowered central densities form \(H_2\) molecules more slowly and are more vulnerable to \(H_2\) photodissociation (Oh & Haiman 2003). The photodissociation is caused by UV photons in the range from 11.6 to 13.6 eV, which overlaps the range from 10.2 to 12.7 eV (from Ly\(\alpha\) to Ly\(\gamma\)) responsible for preheating. This may suppress \(H_2\) cooling and star formation in these minihalos. Minihalos would also be more easily photo-evaporated, thereby reducing their consumption of ionizing photons during reionization (Shapiro et al. 2004).

This preheating also has a strong impact on the redshifted 21 cm signal from the beginning of reionization. UV resonance photons may raise \(T_{H_2}\) above \(T_{\text{CMB}}\), thus moving the gas from absorption to emission (see Fig. 4). In fact, even relatively low intensities (\(<10\) photons per baryon) can drastically affect the signal.

If the first radiation sources deposit a substantial fraction of their energy in X-rays, the average temperature of the IGM may climb even higher (Venkatesan et al. 2001; Oh & Haiman 2003). However, unless the X-ray spectrum is very hard, most of the energy is deposited within ~1 comoving Mpc of the source. By contrast, the UV photons considered here, which travel cosmological distances before scattering, produce a very homogeneous heating mechanism. Therefore, in regions that do not have a radiation source nearby, heating by resonance photons can be more important even if, on average, X-ray heating is stronger.

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REFERENCES

Chen, X., & Miralda-Escude, J. 2004, ApJ, 602, 1
Chuzhoy, L., & Shapiro, P. R. 2005, ApJ, 651, 1
Ciardi, B., & Madau, P. 2003, ApJ, 596, 1
Clementi, E., & Raimondi, D. L. 1963, J. Chem. Phys., 38, 2686
Hirata, C. M. 2006, MNRAS, 367, 259
Madau, P., Meiksin, A., & Rees, M. J. 1997, ApJ, 475, 429

Oh, S. P., & Haiman, Z. 2003, MNRAS, 346, 456
Pritchard, J. R., & Furlanetto, S. R. 2006, MNRAS, 367, 1057
Shapiro, P. R., Iliev, I. T., & Raga, A. C. 2004, MNRAS, 348, 753
Slater, J. C. 1964, J. Chem. Phys., 39, 3199
Venkatesan, A., Giroux, M. L., & Shull, J. M. 2001, ApJ, 563, 1