PROCEEDINGS OF SCIENCE

Dark Matter Annihilation and Hydrogen 21cm Cosmology

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If dark matter is made up of Weakly Interacting Massive Particles, the annihilation of these particles results in partial ionization and heating before the formation of the first stars. We investigate the effect of gas heating on the expected Hydrogen 21cm power spectrum. Heating by particle annihilation results in a minimum in the power spectrum of 21cm fluctuations at the redshift at which the gas temperature equals the CMB temperature. Such a minimum is not expected in the standard theory since the gas temperature in the absence of luminous sources or dark matter heating is always below the CMB temperature. Thus observations of the Hydrogen 21cm power spectrum at multiple redshifts may help us obtain constraints on dark matter particle and halo properties.

International Workshop on Cosmic Structure and Evolution
September 23-25, 2009
Bielefeld , Germany

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1. Introduction.

It is commonly believed that $\sim 80\%$ of the matter in the Universe is dark. Weakly Interacting Massive Particles (WIMPs) are natural dark matter candidates, a common example of which is the supersymmetric neutralino. Neutralinos do not decay, but annihilate in pairs releasing standard model particles like photons, electrons, and neutrinos. Some of these standard model particles interact with gas, resulting in partial ionization and heating. At redshifts $\lesssim 50-60$, we expect particle annihilation to be important mostly in dark matter halos. The earliest WIMPy halos are thought to have formed by redshift $z \sim 60$, with a mass $\sim 10^{-6} M_\odot$ [1]. In this article, we consider WIMP dark matter annihilation and the implications for Hydrogen 21cm cosmology.

Hydrogen 21cm radiation is a valuable probe of the high redshift Universe $z < 200$ [2]. For redshifts $z \lesssim 30$ and before the formation of luminous objects, the Hydrogen 21cm radiation is thought be very small because of the low gas temperature. This scenario may be altered if dark matter annihilation provides a source of heating. Several authors have investigated the effect of particle annihilation on reionization [2, 3, 4, 5] and Hydrogen 21cm spectra [5, 6, 7, 8, 9]. This article is based on our published work [6], to which we refer the interested reader.

2. Particle annihilation in halos.

Let us consider a clump of dark matter particles of mass m_χ, with a density ρ_χ in a volume δV. The probability of 2 WIMPs annihilating in a time δt is $\approx \delta t \langle \sigma_a v \rangle \rho_\chi / m_\chi$. $\langle \sigma_a v \rangle$ is the annihilation cross section times the relative velocity of the particles. The number of particle pairs in the volume δV is $\delta V \rho_\chi / 2m_\chi$. Therefore the number of annihilations per unit volume and per unit time at a redshift z is given by:

$$\frac{dN_{\rm ann}}{dt dV}(z) = \frac{\langle \sigma_{\rm a} v \rangle}{2 m_\chi^2} \rho_\chi^2(z). \tag{2.1}$$

ρ_χ^2 includes contributions from both the halo dark matter component and the free (not in collapsed objects) component. At low redshifts (after halo formation), the halo component is predominant:

$$\rho_\chi^2(z) = (1+z)^3 \int_{\rm M_{min}} dM \, \frac{dn_{\rm halo}}{dM}(M,z) \left[\int_0^{r_{200}} dr \, 4\pi r^2 \rho^2(r) \right] (M,z). \tag{2.2}$$

Before halo formation, we have $\rho_\chi^2(z) = (1+z)^6 \rho_{\rm c}^2 \Omega_{\rm dm}^2$. $\rho_{\rm c}$ is the critical density and $\Omega_{\rm dm}$ is the dark matter fraction today. $n_{\rm halo}$ is the comoving number density of halos and $M_{\rm min}$ is the minimum halo mass. The integral in square brackets is over the halo volume, up to the halo radius r_{200}. Note that the halo luminosity is proportional to the number density of dark matter halos, while the luminosity of free dark matter varies as the *square* of the dark matter density. This is because for halos, the luminosity is mostly due to annihilation of particles within the same halo, and halo-halo collisions are ignored.

For the special case of the NFW density profile [10], the dark matter density within a halo $\rho(r)$ is given by the spherically symmetric form:

$$\rho(r) = \frac{\rho_{\rm s}}{(r/r_{\rm s})\,[1 + r/r_{\rm s}]^2}. \tag{2.3}$$

$\rho_{\rm s}$ and $r_{\rm s}$ are constants. We define the concentration parameter $c_{200} = r_{200}/r_{\rm s}$. r_{200} (which we take to be the halo radius) is the radius at which the enclosed mean density equals $\bar{\rho}$ = 200 $\times$ the cosmological average measured at the redshift $z_{\rm F}(M)$ at which the halo formed. We then have, for NFW halos:

$$\begin{aligned} \int dr 4\pi r^2 \rho^2(r) &= \frac{M\,\bar{\rho}}{3} \left(\frac{\Omega_{\rm dm}}{\Omega_{\rm m}}\right)^2 f[c_{200}(z)] \\ f(x) &= \frac{x^3\,\left[1-(1+x)^{-3}\right]}{3\left[\ln(1+x) - x(1+x)^{-1}\right]^2}. \end{aligned} \tag{2.4}$$

$M = M(r_{200})$ is the mass of the halo, and we have set the mass in dark matter to be $M_{\rm dm} = (\Omega_{\rm dm}/\Omega_{\rm m})\,M$. c_{200} is generally a function of both halo mass and redshift.

Let $dN_\gamma/dE_\gamma(E_\gamma, z)$ be the number of photons per annihilation per photon energy at z. The form of dN_γ/dE_γ depends not only on the particle physics, but also on astrophysical processes such as inverse compton scattering. Inverse compton scattering of CMB photons by energetic ($\sim$ GeV) electrons results in a background of upscattered photons in the $\sim$ MeV range, far more efficient in ionization and heating [4]. As there are $\sim 2\times 10^9$ CMB photons per baryon, we may ignore the mean free path of scattering of high energy electrons from WIMP annihilation with CMB photons (compared to that of high energy photons from WIMP annihilation scattering with gas atoms), and assume that the inverse compton photons are produced approximately at the redshift at which particle annihilation takes place.

We can now write down an expression for the average energy available per gas atom per unit time at a redshift z:

$$\frac{E(z)}{n(z)} = \int_\infty^z \frac{-dz'}{(1+z')H(z')} \left(\frac{1+z}{1+z'}\right)^3 \left(\frac{dN_{\rm ann}}{dtdV}\right)(z') \int_{E_1}^{E_2} dE'_\gamma E'_\gamma \, \frac{dN_\gamma}{dE'_\gamma}(E'_\gamma)\, e^{-\kappa(z',z;E'_\gamma)} \, \left[c\sigma(E'_\gamma)\right]. \tag{2.5}$$

$n(z)$ is the physical gas density at z. We have used the relation $dz = -dt\,(1+z)H(z)$ in Eq. 2.5. The cubic term accounts for the expansion of the Universe. We have set $E'_\gamma = E_\gamma(1+z)/(1+z')$ to account for the redshifting of photon energy, and E_1 and E_2 are appropriately redshifted minimum and maximum values. The exponential term accounts for the scattering of photons as they travel from z' to z. $\sigma(E_\gamma)$ is the scattering cross section for photons of energy E_γ, and c is the speed of light. $\kappa(E_\gamma)$ is the optical depth for photons of energy E_γ, given by the expression

$$\kappa(z', z; E_\gamma) = \int_{z'}^z \frac{-dz''}{(1+z'')H(z'')} \, c\, n(z'')\sigma(E_\gamma). \tag{2.6}$$

Fractions $\eta_{\rm ion}(z)$ and $\eta_{\rm heat}(z)$ of this energy go into ionization and heating respectively [11]. The ionized fraction and gas temperature $T(z)$ are obtained by solving together, the 2 equations:

$$\begin{aligned} (1+z)H(z)\frac{dx_{\rm ion}(z)}{dz} &= -\mu\left[1-x_{\rm ion}(z)\right]\eta_{\rm ion}(z)\left[\frac{E(z)}{n(z)}\right] + n(z)x_{\rm ion}^2(z)\alpha(z) \\ (1+z)H(z)\frac{dT(z)}{dz} &= 2T(z)H(z) - \frac{2\eta_{\rm heat}(z)}{3k_{\rm b}}\left[\frac{E(z)}{n(z)}\right] - \frac{x_{\rm ion}(z)\left[T_\gamma(z)-T(z)\right]}{t_{\rm c}(z)}. \end{aligned} \tag{2.7}$$

μ is the inverse of the average ionization energy per atom, including both H and He. α is the recombination coefficient which depends on $T(z)$. T_γ is the CMB temperature and $t_{\rm c}$ is the compton cooling time scale ≈ 1.44 Myr $[30/(1+z)]^4$. The last term in the temperature evolution equation accounts for the transfer of energy between free electrons and the CMB by compton scattering [7, 12]. Here, we have used $x_{\rm ion} \ll 1$ and have ignored the Helium number fraction.

3. Hydrogen 21cm cosmology.

We now apply the results of the previous section to Hydrogen 21cm cosmology. As the ground state of Hydrogen is hyperfine split into a singlet level and a triplet level, an atom can make a transition from one level to the other, absorbing or emitting a photon of wavelength $\approx$ 21cm. Let us consider a Hydrogen atom at a redshift z. For the wavelengths of interest, we may use the Rayleigh-Jeans approximation, and express the specific intensity of radiation in terms of an equivalent temperature called the brightness temperature $T_{\rm R}$:

$$T_{\rm R}(z) = \left[1 - e^{-\tau(z)}\right] T_{\rm s}(z) + e^{-\tau(z)} T_\gamma(z), \tag{3.1}$$

where τ is the optical depth in 21cm transitions measured to redshift z. $T_{\rm s}$ is the spin (excitation) temperature of the gas. A more useful quantity in 21cm cosmology is the differential brightness temperature defined as $T_{\rm b} = (T_{\rm R} - T_\gamma)/(1+z)$. Solving for $\tau, T_{\rm s}$, and using $\tau \ll 1$, we find [7, 8]:

$$T_{\rm b}(z,\vec{x}) = 27\,{\rm mK}\,\sqrt{\frac{1+z}{10}}\,(1-x_{\rm ion})\,\frac{n}{\bar{n}}\,\frac{\xi}{1+\xi}\left(1 - \frac{T_\gamma}{T}\right)\left[\frac{H(z)/(1+z)}{dv_{||}/dr_{||}}\right]. \tag{3.2}$$

$\bar{n}$ is the spatial average of the gas number density n. ξ is a coupling coefficient which includes both collisional, and Lyman-α couplings. The last term accounts for redshift space distortions due to peculiar velocities. From Eq. 3.2, we see that the differential brightness temperature is sensitive to fluctuations in gas number density ($\delta_{\rm n}$), gas temperature ($\delta_{\rm T}$), and the ionized fraction ($\delta_{\rm x}$). Thus, 21cm observations can teach us much about the physics of the Universe.

The fractional perturbation to the 21cm temperature δ_{21} may be expressed as

$$\delta_{21}(z,\vec{k}) = \delta_{\rm n}(z,\vec{k})\left[\beta_{\rm n}(z) + \left(\hat{n}\cdot\hat{k}\right)^2\right] + \delta_{\rm T}(z,\vec{k})\beta_{\rm T}(z) + \delta_{\rm x}(z,\vec{k})\beta_{\rm x}(z). \tag{3.3}$$

$\beta_{\rm n}$, $\beta_{\rm T}$, and $\beta_{\rm x}$ are functions of z, but not of $\vec{k}$. We expand δ_{21} as a sum over spherical harmonics:

$$\delta_{21}(z,\hat{n}) = \sum_{l,m} a_{lm}(z)\,Y_{lm}(\hat{n}), \tag{3.4}$$

where

$$\begin{aligned} a_{lm}(z) &= \int d\Omega\, Y^*_{lm}(\hat{n}) \int \frac{d^3k}{2\pi^3}\left[\exp -ikr(\hat{k}\cdot\hat{n})\right]\delta_{21}(z,\vec{k}) \\ &= 4\pi(-i)^l \int \frac{d^3k}{(2\pi)^3} Y^*_{lm}(\hat{k})\left[\delta_{\rm n}\left\{\beta_{\rm n} j_l(x) - \frac{\partial^2 j_l(x)}{\partial x^2}\right\} + \delta_{\rm T}\beta_{\rm T} j_l(x) + \delta_{\rm x}\beta_{\rm x} j_l(x)\right]. \end{aligned} \tag{3.5}$$

$x = kr$, and r is the comoving distance

$$r = \int_0^z dz' \frac{c}{H(z')} \approx \frac{2c}{H_0\sqrt{\Omega_{\rm m}}}\left[1 - (1+z)^{-1/2}\right]. \tag{3.6}$$

We can now construct the variance C_l defined as

$$C_l(\nu,\Delta\nu)\,\delta_{ll'}\,\delta_{mm'} = \langle a_{lm}(\nu)\,a^*_{l'm'}(\nu')\rangle. \tag{3.7}$$

$\nu = \nu_0/(1+z)$ is the redshifted frequency, and $\Delta\nu = |\nu-\nu'|$. Here, we consider the simpler case of $\Delta\nu = 0$. For the full solution, see [6]. Let us define the power spectrum

$$\langle\delta_i(z,\vec{k})\delta_j^*(z,\vec{k}')\rangle = (2\pi)^3\delta^3(\vec{k}-\vec{k}')P_{ij}(z,k) = (2\pi)^3\delta^3(\vec{k}-\vec{k}')P_{ji}(z,k), \tag{3.8}$$

where i and j may stand for the baryon number density, ionized fraction, or gas temperature fluctuations.

To compute C_l, we first need to compute the various power spectra in Eq. (3.8). Let us first consider perturbations in the density of dark matter halos. The energy absorbed by the gas at any location z depends on the halo density for redshifts $z' > z$, as seen in Eq. (2.5). The integration over redshifts in Eq. (2.5) results in the perturbations averaging out, as statistically, overdense regions are as likely to exist as underdense regions. As a first approximation, we may ignore perturbations in halo density. The fractional perturbation in E is then $\delta_\mathrm{E} = \delta_\mathrm{n}$. We express the fractional perturbations in the ionized fraction (δ_x) and the gas temperature (δ_T) at redshift z as:

$$\begin{aligned} \delta_\mathrm{x}(z,\vec{k}) &= S_\mathrm{x}(z)\,\delta_\mathrm{n}(z,\vec{k}) \\ \delta_\mathrm{T}(z,\vec{k}) &= S_\mathrm{T}(z)\,\delta_\mathrm{n}(z,\vec{k}), \end{aligned} \tag{3.9}$$

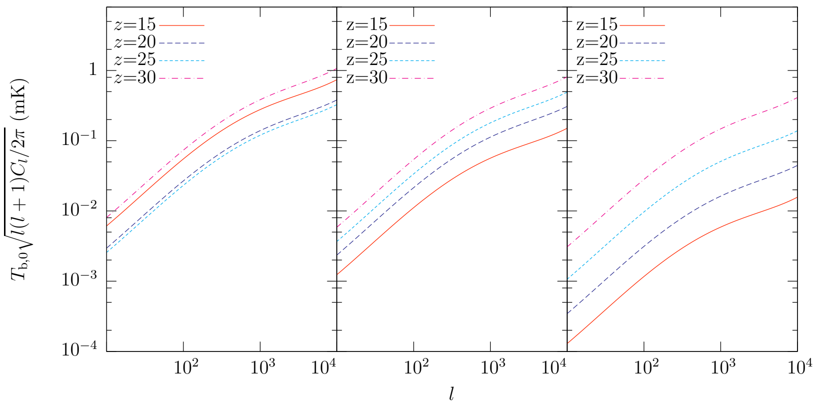
i.e. we make the approximation that the k dependent part of δ_x and δ_T follows the form of δ_n. This is motivated by the dependence of Eq. (2.5), on the baryon density. S_x and S_T give us the time evolution of the perturbations. With this approximation, we consider perturbations in Eq. (2.7) to solve for S_x and S_T. Eq. (3.7) can then be simplified to the form

$$C_l(\nu) = \frac{2}{\pi}\,\frac{1}{(1+z)^2}\int_0^\infty dk\,k^2 P(k)\times\left[\beta(z)\,j_l(x) - \frac{\partial^2 j_l(x)}{\partial x^2}\right]^2, \tag{3.10}$$

where $\beta(z) = \beta_\mathrm{n}(z) + S_\mathrm{T}(z)\beta_\mathrm{T}(z) + S_\mathrm{x}(z)\beta_\mathrm{x}(z)$, and we assumed $P_\mathrm{nn}(z,k) = P(z,k) = P(k)/(1+z)^2$. $P(k)$ is the matter power spectrum given by [13].

4. Results and conclusions.

We now calculate the power spectrum using Eq. 3.10 and Eq. 2.7. Fig. 1 shows the amplitude of fluctuations as a function of multipole l, for different values of redshift z. Fig. 1(a) is drawn for the case of a 10 GeV/c^2 dark matter particle, and halos fitted with an NFW profile with concentration parameter $c_{200} = 5$. Fig. 1(b) corresponds to a 50 GeV/c^2 dark matter particle mass, while 1(c) represents the standard theory with no dark matter heating. Let us compare Fig. 1(a) with 1(c). At a redshift $z = 15$, the amplitude is much larger in (a) than in (c). This is due to the larger gas temperature in model (a) resulting in better coupling between the kinetic temperature and the spin temperature. When z is increased to 20, the amplitude *decreases* in model (a), while it increases in (c). This is because, the gas temperature in model (a) is close to the CMB temperature at this redshift, resulting in a very low value for the differential brightness temperature. In more



(a) $m_\chi = 10$ GeV $c_{200} = 5$ (b) $m_\chi = 50$ GeV $c_{200} = 5$ (c) No DM heating

Figure 1: The power spectrum C_l. Plots (a) and (b) are for dark matter masses $m_\chi = 10$ GeV$/c^2$ and 50 GeV$/c^2$ respectively, with $c_{200} = 5$. (c) is plotted for the standard scenario where heating by dark matter is neglected. The power spectrum shows a minimum with redshift when $T = T_\gamma$ in (a). The minimum is not observed in (b) because $T < T_\gamma$ for all redshifts in the range considered.

physical terms, it is not possible to observe a bright Hydrogen cloud when seen against an equally bright background. In (c), T is always less than T_γ, resulting in a monotonous increase in power with increase in redshift. As the redshift is increased to $z = 30$, the amplitude begins to rise in both (a) and (c). At this redshift, the gas temperature in (a) is below the CMB temperature, and hence there is not much difference between the models (a) and (c). Thus, observing the power spectrum at different redshifts can help us distinguish the dark matter scenario from the standard one. Note however that only dark matter models with small particle mass, large $\langle\sigma_a v\rangle$, or favorable halo parameters show this minimum. No minimum is observed in Fig. 1(b) since the gas temperature in this model is below the CMB temperature for all z in the range considered.

A question arises: Can astrophysical sources mimic heating by dark matter? We first consider Pop. III stars at redshift ~ 25. Pop. III stars are very massive, luminous, and short lived. Each star is surrounded by a hot, ionized region which in turn is enveloped by a warm, neutral region. Using typical values [14], and expected Pop. III star abundances [15], we conclude that Pop. III stars do not heat the Universe as much as light (~ 10 GeV$/c^2$) dark matter, and are not large contaminants. A much larger contribution is made by later astrophysical objects such as quasars and Pop. II stars. [16] have calculated the differential brightness temperature in the presence of astrophysical sources, and find that T_b passes through zero at $z \sim 14-15$. Therefore, it will not be possible to distinguish the dark matter model from the standard theory at redshifts $z \sim 15$ or below. However at redshifts $z > 20$, there is little or no contribution from these baryonic sources. We conclude that observations of the Hydrogen 21cm power spectrum at several redshifts $z \gtrsim 20$ may help us detect the presence of an active dark matter component.

Acknowledgments

We thank Universität Bielefeld and the Deutsche Forschungsgemeinschaft for financial support.

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