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Sunyaev–Zeldovich fluctuations from the first stars?

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ABSTRACT

WMAP's detection of high electron-scattering optical depth τ_e suggests substantial star formation at high redshift $z \sim 17 \pm 5$. On the other hand, the recovered $\sigma_8 \sim 0.84 \pm 0.04$ argues against a cluster Sunyaev–Zeldovich (SZ) origin for the observed small-scale cosmic microwave background (CMB) fluctuation excess, which generally requires $\sigma_8 \sim 1.1$. Here we consider the effects of high-redshift star formation on the CMB. We derive a fairly model-independent relation between τ_e and the number of ionizing photons emitted per baryon N_γ, and use this to calibrate the amount of high-redshift supernova activity. The resulting supernova remnants Compton cool against the CMB creating a Compton y distortion $y \sim$ few $\times 10^{-6}$ within observational bounds. However they also create small-scale SZ fluctuations, which could be comparable with SZ fluctuations from unresolved galaxy clusters. This raises the exciting possibility that we have already detected signatures of the first stars not just once, but twice, in the CMB.

Key words: galaxies: formation – intergalactic medium – cosmic microwave background.

1 INTRODUCTION

The recent detection of high electron-scattering optical depth $\tau = 0.17 \pm 0.04$ by the Wilkinson Microwave Anisotropy Probe (WMAP) suggests a reionization redshift $z_r = 17 \pm 5$ (Kogut et al. 2003; Spergel et al. 2003), providing good evidence for significant star formation (SF) at high redshift z. WMAP combined with other large-scale structure data also supports a ΛCDM cosmology with power-spectrum normalization $\sigma_8 = 0.84 \pm 0.04$.

This power-spectrum normalization is discrepant from that inferred from the cosmic microwave background (CMB) fluctuation excess at small scales (Dawson et al. 2002; Mason et al. 2003), if this excess is attributed to the Sunyaev–Zeldovich (SZ) effect from unresolved groups and clusters (Bond et al. 2002; Komatsu & Seljak 2002; Goldstein et al. 2002). These observations require $\sigma_8(\Omega_b h/0.035)^{0.29} = 1.04 \pm 0.12$ at the 95 per cent confidence level (Komatsu & Seljak 2002).

It has been argued that galactic winds could give rise to a detectable SZ effect (Majumdar, Nath & Chiba 2001). Here we argue that the stellar activity required to photoionize the Universe at $z_r \sim 20$ injects a considerable amount of energy into the intergalactic medium (IGM), which is then transferred to the CMB owing to the efficiency of Compton cooling at these high redshifts. Although the resulting mean Compton-y distortion is consistent with the experimental upper limit, there may be detectable angular fluctuations in the y distortion. We show, in fact, that for reasonable reionization parameters the fluctuation amplitude from high-z SF may be comparable to that from galaxy clusters. If so, then the above-mentioned discrepancy in the power-spectrum normalization may be resolved.

In the next section we argue that supernova remnants at $z \gtrsim 10$ cool by Compton heating of the CMB and discuss the energetics of this process. In Section 3, we derive a relation between the measured optical depth τ_e and the number of ionizing photons required to reionize the Universe. We then show that this number of ionizing photons is proportional to the energy injected into the IGM by supernovae, and thus the energy transferred to the CMB. In Section 4 we discuss angular fluctuations in the y distortion and show that they may be comparable at small scales to those from unresolved clusters.

In all numerical estimates, we assume a Λ CDM cosmology given by the best fits to the WMAP data: $(\Omega_m, \Omega_\Lambda, \Omega_b, h, \sigma_8) = (0.27, 0.73, 0.044, 0.7, 0.84)$.

2 HOW DOES THE SUPERNOVA REMNANT COOL?

At redshifts $z > 7$, galactic winds powered by multiple ($>10^5$) supernovae (SN) or an energetic quasar jet are cooled primarily by Compton cooling from the CMB (Tegmark, Silk & Evrard 1993; Voit 1996; Madau, Ferrara & Rees 2001). Less powerful winds result in cooler remnants where radiative losses could potentially be important. However, at $z \sim 10$–20, the wind from even a single SN will lose a substantial fraction of its energy to the CMB, as we show below.

Zero-metallicity stars should be supermassive, $M_* \geqslant 100\ M_\odot$, owing to the thermodynamics of molecular hydrogen (H_2) cooling (Abel, Bryan & Norman 2000; Bromm, Coppi & Larson 2002).

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Furthermore, pair-instability SN from such very massive stars (VMSs) should have explosion energies ~100 times more powerful than conventional type II SN, $E_{\rm VMS} \sim 10^{53}$ erg (Heger & Woosley 2002). An extreme but plausible version of zero-metallicity SF in low-mass halos $T_{\rm vir} < 10^4$ K prevalent at high z is 'one star per halo', where internal UV photodissociation of H_2 by the first star in that halo halts all further gas cooling and SF (Omukai & Nishi 1999; Glover & Brandt 2001). Simulations show the VMS quickly photoevaporates all the gas within the shallow halo potential well within a sound-crossing time (M. Norman, private communication). In addition, the VMS photoionizes a region around the halo up to $R \sim 70\ (M_{\rm VMS}/100\ {\rm M}_\odot)^{1/3}$ kpc comoving, assuming that each baryon in the VMS can ionize $\sim 10^5$ H I atoms (Bromm, Kudrizki & Loeb 2001). Thus, the SN remnant (SNR) expands into a pre-ionized region at roughly the mean IGM gas density. During the adiabatic Sedov phase, $R = \gamma_{\rm o}(Et^2/\rho_{\rm IGM})^{1/5}$, where $\gamma_{\rm o} = 1.17$. The remnant is no longer adiabatic and begins to Compton cool when $t \approx t_{\rm C}$, where the Compton cooling time is

$$t_{\rm C} = 3 m_{\rm e} c \left(4\sigma_T a T_{\rm CMB}^4\right)^{-1} = 1.4 \times 10^7 [(1+z)/20]^{-4}\ {\rm yr}, \qquad (1)$$

independent of temperature and density. The (proper) size of the remnant at this point, when it quickly loses most of its energy, is

$$R = 2.2 \left(E_{\rm VMS}/10^{53}\ {\rm erg}\right)^{1/5} [(1+z)/20]^{-11/5}\ {\rm kpc} \qquad (2)$$

in physical units. The angular scale is $\theta = R/d_A = 0.9$ arcsec $(E_{\rm VMS}/10^{53}\ {\rm erg})^{1/5}[(1+z)/20]^{-11/5}$ (which corresponds to $l = \pi/\theta = 7.6 \times 10^5$), beyond the reach of present-day CMB interferometers. Thus, SNRs are effectively point sources, unless many SN explode together in the same galaxy, and/or SN bubbles from clustered haloes overlap (see below).

Most of the mass and energy of the remnant is in the dense post-shock shell, which is at $\rho_{\rm shell} \sim 4\ \rho_{\rm IGM}$. At $t = t_{\rm C}$, we can compute the temperature behind the shock front from the Sedov–Taylor solution, $v_s = 0.4\gamma_{\rm o}(E/\rho_{\rm IGM} t^3)^{1/5}$, and assuming a strong shock $T_{\rm c} = 3 v_{\rm s}^2 \mu m_{\rm p}/16 k_{\rm B}$. We thus obtain the ratio of Compton and isobaric radiative cooling time $t_{\rm rad} = 2.5\ k_{\rm B} T/[n\Lambda(T)]$ at $t = t_{\rm C}$ as

$$\frac{t_{\rm rad}}{t_{\rm C}} = 0.4 \left(\frac{E_{\rm VMS}}{10^{53}\ {\rm erg}}\right)^{0.4} \left(\frac{1+z}{20}\right)^{4.6} \Lambda_{23}^{-1}, \qquad (3)$$

where $\Lambda(T) = \Lambda_{23} \times 10^{-23}$ erg s^{-1} cm^3, and $\Lambda_{23} \sim 1$ for low-metallicity gas with $T \sim 10^5$–10^7 K. Thus, roughly a third of the SNR energy is lost to Compton cooling.

The electron–ion equilibration time-scale

$$t_{\rm ei} = 10^5 {\rm yr} \left(\frac{1+z}{15}\right)^{-3} \left(\frac{\delta}{4}\right)^{-1} \left(\frac{T}{10^6\ {\rm K}}\right)^{3/2} \qquad (4)$$

(where δ is the overdensity of the post-shock shell) is significantly shorter than the Compton cooling time at all redshifts, so there is no problem in quickly transferring the shock energy from protons to electrons.

A scenario that is perhaps more likely is one where many stars $M_{*,\rm tot} \sim 10^7 (f_*/0.1)\ (f_{\rm b}/0.1)\ (M_{\rm DM}/10^9)\ {\rm M}_\odot$ (where $f_{\rm b} \approx \Omega_{\rm b}/\Omega_{\rm m}$ is the baryon fraction, and f_* is the fraction of baryons that fragment to form stars) form together in rarer, more massive haloes $T_{\rm vir} > 10^4$ K, where atomic cooling allows much higher gas densities and more efficient SF (Oh & Haiman 2002). The massive-star evolution time-scale is $t_* \sim 3 \times 10^6\ {\rm yr} \ll t_{\rm C}$. Thus, if SF takes place in a starburst mode, the explosions are essentially simultaneous, and $E_{\rm tot} \approx N_{\rm SN} E_{\rm SN}$. Then, an extremely energetic wind powers a much hotter bubble, and from equation (3), $t_{\rm rad}/t_{\rm C} \propto E_{\rm tot}^{0.4} \gg 1$ and radiative cooling is entirely negligible. For instance, if $f_* \sim 10$ per cent of the baryons in a $M_{\rm DM} \sim 10^9\ {\rm M}_\odot$ halo fragment to form VMSs, $t_{\rm C} \sim t_{\rm rad}/40$). In principle, radiative losses could be significant in the dense interstellar medium (ISM) of these larger haloes (since photoevaporation does not take place in these deeper potential wells); however in practice most simulations (e.g. Mac-Low & Ferrara 1999) find that for such low-mass systems, the SN bubbles quickly 'blow out' (particularly in discs) and vent most of their energy and hot gas into the surrounding IGM. Hereafter we shall encapsulate this uncertainty as $\epsilon \approx 0.3$–1, the average fraction of the explosion energy lost to the CMB via Compton cooling. If stars form in clusters in higher-mass haloes rather than singly in low-mass haloes we expect this efficiency to be high, $\epsilon \geqslant 0.8$.

The spatial distribution of SF does not affect our estimate of the mean Compton-y distortion ($y = k_{\rm B} T_{\rm hot}/(m_{\rm e} c^2)\tau_{\rm hot}$): more clustered SF results in higher $T_{\rm e}$ but lower $\tau_{\rm hot}$. However, it does of course affect the strength of SZ fluctuations. We now turn to these issues.

3 THERMAL SUNYAEV–ZELDOVICH EFFECTS

3.1 SZ flux from individual supernovae

The SZ flux from an individual SNR is

$$\begin{aligned} S_\nu &= \frac{2k_{\rm B}^3 T_\gamma^2}{h^2 c^2} g(x) \int {\rm d}\Omega |\Delta T_\nu(\theta)| \\ &= \frac{2k_{\rm B}^3 T_\gamma^3}{h^2 c^2} g(x) \frac{k_{\rm B} T_{\rm e}}{m_{\rm e} c^2} \sigma_{\rm T} \frac{N_{\rm e}}{d_A^2} \\ &= 1.8 \times 10^{-2} \left(\frac{g(x)}{4}\right) \left(\frac{E_{\rm VMS}}{10^{53}\ {\rm erg}}\right) \left(\frac{\epsilon}{0.5}\right) \left(\frac{z}{20}\right)^2 {\rm nJy}, \end{aligned} \qquad (5)$$

where $g(x) = x^4 {\rm e}^x [x \coth(x/2) - 4]/({\rm e}^x - 1)^2$ is the spectral function, $x \equiv h\nu/kT_\gamma$, $T_\gamma = 2.7$ K is the CMB temperature, and $N_{\rm e}$ is the total number of hot electrons at temperature $T_{\rm e}$. In the second line we have used $K_{\rm b} T_{\rm e} N_{\rm e} \approx E_{\rm VMS}$. The energy of the remnant is a function of time, $E_{\rm VMS}(t) \approx E_{\rm VMS,o} \exp(-t/t_{\rm C})$ (in the regime where Compton cooling off the CMB dominates). The flux from an individual SNR is well beyond threshold for any realistic experiment; only a very large number of SN ($>10^8$) going off simultaneously within a star cluster will be detectable at the ~ mJy level. Thus, SN bubbles cannot be identified and removed from SZ maps; unresolved SN will create both a mean Compton-y distortion and temperature fluctuations, which we now calculate.

3.2 Mean Compton y distortion

We first use the observed optical depth $\tau_{\rm e}$ to derive a lower limit to the number of ionizing photons N_γ emitted per baryon. The dominant contribution to $\tau \propto (1+z)^{1.5}$ comes from high z where the recombination time $t_{\rm rec} \propto (1+z)^{-3}$ is short, and recombinations are the rate-limiting step toward achieving reionization. The filling factor of H II regions is $Q_{\rm H\,II} \approx t_{\rm rec}/t_{\rm ion} \approx \dot{N}_\gamma/[\alpha_{\rm B} n_{\rm e}(z) C_{\rm II}(z)]$, where $\dot{N}_\gamma$ is the rate at which ionizing photons are emitted per baryon (in units of s^{-1}), and $C_{\rm II} \equiv \langle n_{\rm e}^2\rangle/\langle n_{\rm e}\rangle^2$ is the clumping factor of ionized regions (e.g. Madau, Haardt & Rees 1999). The clumping factor increases with time as structure formation proceeds; it declines sharply at high z and is $C \approx 2$ at $z = 20$, compared to $C \approx 30$ at $z = 10$ (Haiman, Abel & Madau 2001). More sophisticated considerations (Miralda-Escudé et al. 2000) take into account the density dependence of reionization, but apply primarily near the epoch of overlap, $Q_{\rm II} \to 1$, when overdense regions are ionized. This has little impact on our estimates. Most of the mass and the optical depth comes from regions close to the mean density.

The electron-scattering optical depth is given by

$$\tau_{\rm e} = c\sigma_{\rm T} \int {\rm d}z \frac{{\rm d}t}{{\rm d}z} n_{\rm e}(z) {\rm min}\left(1, \frac{\dot{N}_\gamma}{\alpha_{\rm B} n_{\rm e}(z) C_{\rm II}}\right) = \frac{c\sigma_{\rm T} N_\gamma}{\alpha_{\rm B} C_{\rm II}}. \quad (6)$$

Owing to the cancellation of the electron density, this expression is *independent* of the redshift of reionization, and the evolution of the comoving emissivity $\dot{N}_\gamma(z)$ with redshift; it allows us to relate $\tau_{\rm e}$ and N_γ directly. The only redshift dependence lies in the effective clumping factor $C_{\rm II}$, which increases if reionization takes place at late times. The second equality breaks down if overlap $Q_{\rm II} \rightarrow 1$ is achieved at high z and $\dot{N}_\gamma/[\alpha_{\rm B} n_{\rm e}(z) C_{\rm II}(z)] > 1$ (i.e. recombinations no longer balance ionizations); in using the expression we would then underestimate N_γ, which would only imply an even larger emissivity. The high optical depth $\tau_{\rm e} = 0.17 \pm 0.08$ (2σ) (Kogut et al. 2003; Spergel et al. 2003) detected by WMAP therefore implies that

$$N_\gamma^{\rm IGM} = 17 \pm 8(\bar{T}/10^4)^{-0.7}\,(C_{\rm II}/4) \quad (7)$$

ionizing photons were emitted per baryon, where $\bar{T}$ is the mass-weighted temperature of the reionized IGM (the $\bar{T}^{-0.7}$ factor arises from the temperature dependence of the recombination coefficient). Consistency with WMAP requires more SF if reionization took place at lower redshift, due to the increase in gas clumping at late times.

Since only a fraction $f_{\rm esc}$ of ionizing photons escape from their host halo owing to photoelectric absorption, the actual total number of ionizing photons produced is larger, $N_\gamma^{\rm tot} = N_\gamma^{\rm IGM} f_{\rm esc}^{-1}$. In addition, we only care about those photons emitted at $z > 6$, when $t_{\rm C} < t_{\rm H}$ (where $t_{\rm H}$ is the Hubble time) and Compton cooling is most efficient. Since $\tau_{\rm e}(z < 6) \approx 0.05$, we have $\tau_{\rm e}(z > 6) \approx 0.12$; therefore $0.12/0.17 \sim 0.7$ of the photons are emitted at $z > 6$. Thus

$$N_\gamma^{\rm tot}(z > 6) \approx 25(f_{\rm esc}/0.3)^{-1}\,[\tau_{\rm e}(z > 6)/0.12]\,(C_{\rm II}/4)\,. \quad (8)$$

Estimates for the escape fraction span $f_{\rm esc} \sim 10^{-2} - 1$, but if the earliest stars reside in low-mass haloes with $T_{\rm vir} < 10^4$ K, the gas in such haloes is quickly photoionized and driven out in a photoevaporating wind. If so, $f_{\rm esc} \sim$ few $\times$ 0.1 to $f_{\rm esc} \sim 1$.

How much SF and energy production is associated with $N_\gamma^{\rm tot}$? We consider first VMSs, supported as the source of reionization perhaps by elemental-abundance evidence from low-metallicity halo stars (Oh et al. 2001) and theoretical modelling (Cen 2002; Wyithe & Loeb 2003). Bromm et al. (2001) find that for 300 $M_\odot < M_* < 1000 M_*$, the luminosity per solar mass is approximately constant; for $M_* \sim 100\ M_\odot$, it falls by a factor of 2. Our estimates are thus independent of initial mass function (IMF) details. For one ionizing photon per baryon in the Universe, $f_* \sim 10^{-5}$ baryons have to be processed into VMSs; thus, $N_\gamma^{\rm tot} = 25$ corresponds to $f_* \sim 2.5 \times 10^{-4}$. A $\sim$100–$M_\odot$ pair-instability SN releases $E_{\rm VMS} \sim 10^{53}$ erg (Heger & Woosley 2002), or $E_{\rm b} \sim 0.5$ MeV per baryon processed into the VMS. The total energy release per baryon is therefore

$$E_{\rm c} = \epsilon f_* E_{\rm b} = 100\,(\epsilon/0.8)\left(N_\gamma^{\rm tot}/25\right)\,{\rm eV}, \quad (9)$$

where ϵ is the fraction of the thermal energy that is lost to the CMB. A possible caveat is the case where a large fraction of the mass in the first stars went into VMSs with $M_* > 260\ M_\odot$, which may collapse directly to black holes without exploding as SN (Heger & Woosley 2002).

The fraction of baryons processed into VMSs $f_* \sim 2.5 \times 10^{-4}(N_\gamma/25)$ implies an IGM metallicity $Z \sim 6 \times 10^{-3}\ Z_\odot$, assuming uniform enrichment (since $\sim$ half the VMS mass is thought to end up as metals). This is consistent with the observed metallicity of the Lyα forest at $z = 3$ of $Z \approx 10^{-2.5}\ Z_\odot$, which is not observed to evolve strongly at higher z (Songaila 2001). Thus, the metals seen in the Lyα forest may well have been injected at very high z by Pop III stars. No trace of the entropy injection associated with the metal-polluting winds would remain, owing to the high efficiency of Compton cooling.

Our derived ionizing-photon:energy:metal ratios would also hold for normal stellar populations (rather than VMSs), which produce roughly the same amount of SN energy and metals per ionizing photon. The arguments are also roughly independent of IMF, as the massive stars that emit ionizing photons also eventually explode as SN.

We now compute the Compton-y parameter associated with this energy injection. For simplicity, we assume that all of the energy is injected at some redshift z_i. The actual redshift evolution introduces at most a factor $\sim$2 uncertainty (see expression below). The y parameter is then given by

$$\begin{aligned} y &= \frac{c\sigma_{\rm T}}{m_{\rm e}c^2}\int_{t_{\rm i}}^{t_0} {\rm d}t\, n_{\rm e}(t) E_{\rm c,o} {\rm e}^{-(t-t_{\rm i})/t_{\rm C}} \\ &\approx n_{\rm e}(z_{\rm i})\sigma_{\rm T} c t_{\rm C}(z_{\rm i}) \frac{E_{\rm c}}{m_{\rm e}c^2} \\ &= 3.6 \times 10^{-6}(1+z_{\rm i}/15)^{-1}(E_{\rm c}/100\,{\rm eV}), \end{aligned} \quad (10)$$

where we have moved the electron density outside the integrand, $n_{\rm e}(t) \approx$ const, since the density does not change significantly on the time-scale over which the gas Compton cools. In the RJ limit, $(\Delta T/T) = -2y = 7 \times 10^{-6}$. The y distortion is less than the *COBE* FIRAS constraint, $y \leqslant 1.5 \times 10^{-5}$ (Fixsen et al. 1996), as it should be. Such a y distortion could in principle be detected by future instruments (Fixsen & Mather 2002). In addition, a low-frequency distortion arising from free–free emission from ionized haloes should also be present (Oh 1999), which should be detectable by future missions such as the *Diffuse Microwave Emission Survey* (*DIMES*).[1]

We pause here for a simple order-of-magnitude check. Let the total amount of energy per baryon injected through Compton cooling into the CMB be $E_{\rm c}$. If this takes place at some median redshift $z_{\rm i}$, this introduces an energy density perturbation of the CMB $\Delta U_\gamma \sim n_{\rm b} E_{\rm c} \sim 6.8 \times 10^{-2}\,([1+z]/15)^3\,(E_{\rm c}/100\,{\rm eV})\,{\rm eV\,cm^{-3}}$. The CMB energy density is $U_\gamma = 1.3 \times 10^4[(1+z)/15]^4\,{\rm eV\,cm^{-3}}$, resulting in a temperature perturbation

$$\frac{\Delta T}{T_\gamma} \sim \frac{1}{4}\frac{\Delta U}{U} \sim 5.2 \times 10^{-6}\left(\frac{1+z}{15}\right)^{-1}\left(\frac{E_{\rm c}}{100\,{\rm eV}}\right) \quad (11)$$

roughly consistent with our previous estimate, from $(\Delta T/T) = -2y$. Why is the mean y distortion arising from non-gravitational heating by high-z SN competitive with that from galaxy clusters today? By integrating over the Press–Schechter mass function and assuming $T_{\rm gas} = T_{\rm vir}$, we find that the mean mass-weighted gas temperature today is $\langle T \rangle = 0.7$ keV. However the Compton cooling time in clusters is $t_{\rm C} \sim 150 t_{\rm H}$, so only $\epsilon \sim 1/150$ of that energy is extracted. Since $y \propto E_{\rm c}(1+z)^{-1}$, we find that the y distortion owing to clusters is $\sim$(0.7 keV/150)/0.1 keV $\times$ 15 $\sim$ 1 times the distortion arising from high-z SN.

4 SZ FLUCTUATIONS

We now calculate the CMB fluctuations induced by high-z SN. We suppose for simplicity that stars form only in haloes where atomic cooling can operate, $T_{\rm vir} > 10^4$ K, where some constant

[1] see http://map.gsfc.nasa.gov/DIMES/

fraction f_* of the baryonic mass fragments to form stars. We use Press–Schechter theory to calculate the abundance of haloes. A hot bubble around each source has a total flux $\propto E_{\rm SN}$ as given by equation (5), and lasts for a Compton cooling time $t_{\rm C}$. The size of the hot bubble is given by equation (2). The finite bubble size damps the power spectrum on scales below the bubble size. For simplicity we shall assume $y_l = y_{\rm o} \exp\,[-(l/l_{\rm c})^2]$, where y_l is the Fourier transform of the y profile of the bubble, $l_{\rm c} = \pi/\theta_{\rm c}$ and $\theta_{\rm c}$ is the angular size of the bubble when most of the Compton cooling takes place. If the SF efficiency is independent of halo mass then $y_{\rm o} = RM_{\rm halo}$, where the normalization constant $R \propto \epsilon f_*$ is determined from the condition that

$$\bar{y} = \int {\rm d}z({\rm d}V/{\rm d}z\,{\rm d}\Omega) \int_{M_{\rm min}}^{\infty} {\rm d}M({\rm d}n/{\rm d}M) y_{\rm o}(M, z), \qquad (12)$$

subject of course to the condition that $f_{\rm coll} > f_{\rm VMS}$ and $\epsilon\ f_* < 1$.

In reionized regions, gas accretion is suppressed in haloes with $T_{\rm vir} < T_{\rm min} \approx 2.5 \times 10^5$ K (or $v_{\rm cir} \sim 50$ km s^{-1} Thoul & Weinberg 1996); lower-mass haloes are thus unlikely to be able to form stars. This boosts the clustering bias of SF systems as reionization proceeds, which increases the strength of SZ anisotropies. It also increases Poisson fluctuations since such massive haloes are rarer. To keep our analysis general, we conservatively only require $T_{\rm vir} > 10^4$ K, but then show how increasing the Jeans mass would boost CMB fluctuations.

The Compton y power spectrum arising from clustering of sources is given by

$$C_l(y) = \int {\rm d}z \frac{{\rm d}V}{{\rm d}z{\rm d}\Omega} P\left(k = \frac{l}{d_M(z)}\right) \times \left[\int_{M_{\rm min}}^{\infty} {\rm d}M \frac{{\rm d}n}{{\rm d}M} D(z) b(M, z) y_l(M, z)\right]^2 \qquad (13)$$

where $P(k)$ is the linear power spectrum, $d_M = d_A(1+z)$ is the comoving angular diameter distance, $D(z)$ is the linear theory growth factor, and $b(M, z)$ is the linear bias factor (Mo & White 1996). We have used the Limber approximation $k = l/d_M$ which is valid for small angles. Note that $C_l(\Delta T/T) = 4\ C_l(y)$ in the RJ limit. The results are shown in Fig. 1 for three cases: (A) a maximal case with $y = 10^{-5}$ (consistent with the current uncertainty in $\tau_{\rm e}$, $C_{\rm II}$, and $f_{\rm esc}$), the largest value allowed by the *COBE* constraint $y < 1.5 \times 10^{-5}$), and clustering bias associated with $T_{\rm vir} > 10^5$ K haloes; (B) a standard 'best-estimate' case with $y = 3.6 \times 10^{-6}$ and clustering bias associated with $T_{\rm vir} > 2 \times 10^4$ K haloes; (C) a minimal scenario with $y = 10^{-6}$ and clustering bias associated with $T_{\rm vir} >$ 5000 K haloes. Also shown are the cluster-induced power spectra for $\sigma_8 = 0.84 \pm 0.08$ (2σ), computed as in Cooray (2000). Although the 'best-estimate' reionization signal lies below the cluster signal, with current uncertainties they could plausibly be comparable. The shapes of the power spectra are fairly well constrained, but their amplitudes are uncertain by ~2 orders of magnitude, as we discuss below.

Roughly speaking, the CMB power spectrum is $C_l \approx \bar{y}^2 w_l$, where $\bar{y}$ is the mean Compton y parameter from equation (11), and $w_l \propto l^n$ [if $P(k) \propto k^n$] is the flux-weighted halo angular power spectrum. The flatness of $w_l l^2$ at high l is because $P(k) \propto k^{-2}$ at these wavenumbers. For haloes with $T_{\rm vir} > 10^4$ K, the rapid increase in bias tend to cancel the decrease in the growth factor at high z, and the halo correlation function and power spectrum $b(M_1)b(M_2)D(z)^2 P(k)$ do not evolve strongly with redshift. We see this in Fig. 2, where we plot $[b(M(\tilde{T}_{\rm c}), z)D(z)]^2$, and $\tilde{b}(M(T_{\rm c}), z)$ is the mass-weighted bias,

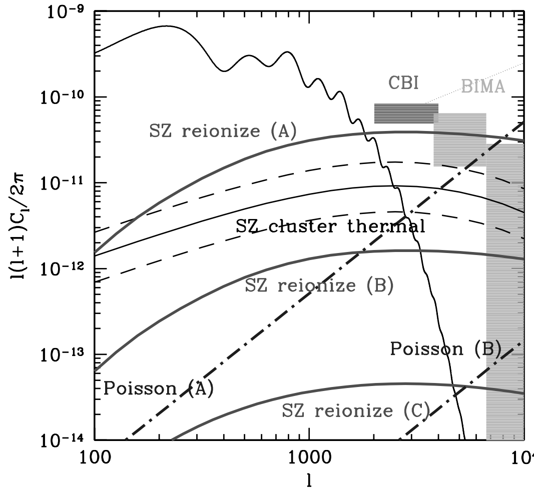

Figure 1. The small-scale power spectrum of the CMB. Our standard case B for thermal SZ from reionization arising from clustering assumes $y = 3.6 \times 10^{-6}$ and clustering bias associated with $T_{\rm vir} > 2 \times 10^4$ K haloes, while a maximal case A assumes $y = 10^{-5}$, and clustering bias associated with $T_{\rm vir} > 10^5$ K haloes. A minimal case C assumes $y = 10^{-6}$ and clustering bias associated with $T_{\rm vir} > 5000$ K haloes. At fixed $\bar{y}$ and halo mass threshold, the Poisson spectra (dot–dashed lines) are more uncertain than the clustering signal; they lie below the clustering signal even for the fairly extreme case shown here (see text). Also shown are the power spectra arising from physics at the surface of last scatter, and the thermal SZ effect for clusters for $\sigma_8 = 0.84 \pm 0.08$ (2σ) (dashed lines are curves for $\sigma_8 = 0.76, 0.92$).

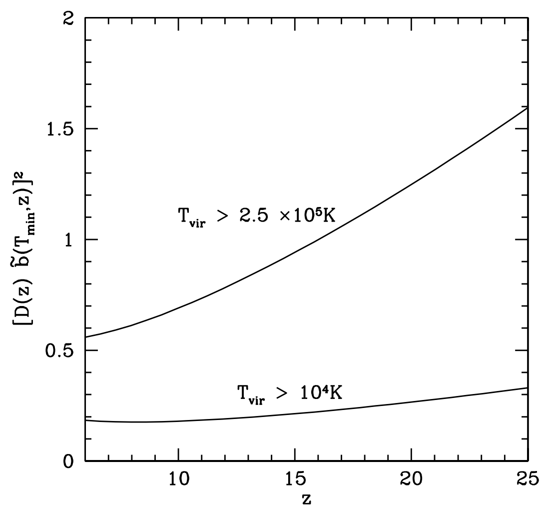

Figure 2. The evolution of the normalization factor for the halo power spectrum, $[D(z)\tilde{b}(M, z)]^2$, where $D(z)$ is the growth factor and $\tilde{b}(M, z)$ is the mass-weighted bias from equation (14). The halo power spectrum for haloes where $T_{\rm vir} > 10^4$ K (which can cool by atomic cooling) hardly evolves with redshift; the increase in bias cancels the decrease in the growth factor. In addition, as reionization proceeds, haloes with $T_{\rm vir} < 2.5 \times 10^5$ K are unable to accrete gas; hence the clustering bias of SF haloes will increase further. The thermal SZ power spectra in Fig. 1 scale directly with this factor.

$$\tilde{b}(M_{\rm c}, z) = \int_{M_{\rm c}}^{\infty} {\rm d}M \frac{{\rm d}N}{{\rm d}M} M b(M, z) \Bigg/ \int_{M_{\rm c}}^{\infty} {\rm d}M \frac{{\rm d}N}{{\rm d}M} M, \qquad (14)$$

which corresponds to the flux-weighted bias since we assume $S \propto M$. This is likely a minimal estimate of the bias since the SF efficiency (and hence the thermal SZ flux) is likely to increase with the depth of the potential well. As reionization proceeds, the actual bias interpolates between the two curves, since accretion is suppressed in haloes forming in reionized regions with $T_{\rm vir} < 2.5 \times 10^5$ K; it approaches the upper curve as $Q_{\rm II} \rightarrow 1$. Since we are probing scales on order of or smaller than the halo correlation length, $r_{\rm o} \sim$ few Mpc comoving, it is reasonable to expect projected halo density (and hence flux) enhancements of order $l\,(l+1)w_l/(2\pi) \sim$ few.

The Poisson power spectrum is given by

$$C_l^{\rm Poisson} = \int {\rm d}z \frac{{\rm d}V}{{\rm d}z{\rm d}\Omega} \int_{M_{\rm min}}^{\infty} {\rm d}M \frac{{\rm d}n}{{\rm d}M} |y_l(M,z)|^2. \quad (15)$$

Because it depends on the detailed shape of the number counts, the Poisson signal is more uncertain than the clustering signal. In particular, the Poisson signal is dominated by the rarest, brightest sources, and is much more sensitive to our uncertain star formation prescription. In Fig. 1 we place an upper bound on Poisson fluctuations by assuming that the sources responsible for a given $\bar{y}$ lie at $z > 12$, when the collapsed mass fraction is small; this implies that very rare bright sources dominate (dot–dashed lines). For case (A), this translates into $f_* \sim 100$ per cent of the baryons in haloes with $T > 10^5$ K fragmenting to form stars. The Poisson signal is significantly less than the clustering signal except for extremely small angular scales $l > 10^4$. This is in contrast to the galaxy cluster SZ signal, when Poisson fluctuations dominate. This is easy to understand: like high-z haloes, clusters are $\sim$2–3σ fluctuations at the epoch at which they form and contain roughly the same fraction of collapsed mass; however, they are more massive by $\sim$6 orders of magnitude and hence have a much lower space density. Each cluster contributes a much larger fraction of the total background signal, increasing the strength of Poisson fluctuations.

Overall, our primary uncertainties in the predicted amplitude are due to uncertainties in the mean y parameter that arise from the uncertainties in $\tau_{\rm e}$, $C_{\rm II}$ and $f_{\rm esc}$ discussed above. There is then an additional uncertainty of $\sim$ few introduced by the range of halo bias factors illustrated in Fig. 2.

5 CONCLUSIONS

We have pointed out that WMAP's large electron-scattering optical depth $\tau_{\rm e}$ implies that SZ fluctuations from high-z SF could be considerable. As an interesting secondary result, we derive a relation between $\tau_{\rm e}$ and N_γ, the number of ionizing photons emitted per baryon. We use this to calibrate the amount of SN activity, and thereby obtain the expected Compton-y distortion, $y \sim$ few $\times 10^{-6}$. Fluctuations in the Compton-y parameter could be detectable and may well account for the the small-scale CMB fluctuation excess at small angular scales. If so, small-scale CMB measurements are *not* a reliable independent measure of σ_8. If the small-scale CMB anisotropies are caused by clusters alone, they will be resolved by forthcoming high-sensitivity and high-resolution SZ surveys. On the other hand, if high-z SF contributes significantly, there will be a substantial unresolved component, since the extremely faint flux from individual haloes is undetectable. A large amount of high redshift SN activity also produces X-rays (Oh 2001), with interesting consequences for reionization.

If a high-z origin of the observed small-scale CMB fluctuations is confirmed, CMB maps may then be used to study the topology of reionization, perhaps by cross-correlating with future 21-cm tomographic maps of neutral hydrogen at high z (Tozzi et al. 2000). Here we have focused exclusively on thermal-SZ fluctuations, which induce a Compton-y distortion to the CMB frequency spectrum and can thus be distinguished from 'genuine' temperature fluctuations with multifrequency CMB measurements. However, high-z SF may also induce temperature fluctuations by scattering from reionized regions with coherent large-scale peculiar velocities, as we detail in a forthcoming paper (Cooray et al., in preparation).

Given the uncertainties in high-z SF discussed above, we can make predictions for small-scale y fluctuations with roughly an order-of-magnitude level of uncertainty in the CMB fluctuation amplitude, and thus cannot at this point conclusively attribute observed small-scale CMB fluctuation excesses to high-z star formation. Nonetheless, this interpretation of the excess is certainly plausible. If it is correct, then the CMB experimentalists have achieved a remarkable triumph: not only have they fulfilled a decade-old quest to measure cosmological parameters with exquisite and unprecedented precision, they have detected signatures of the very first generation of star formation not just once, but twice.

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