Impact of primordial ultracompact minihaloes on the intergalactic medium and first structure formation

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ABSTRACT

The effects of dark matter annihilation on the evolution of the intergalactic medium (IGM) in the early Universe will be more important if the dark matter structure is more concentrated. Ultracompact minihaloes (UCMHs), which formed through dark matter accretion on to primordial black holes (PBHs) or an initial dark matter overdensity produced by a primordial density perturbation, provide a new type of compact dark matter structure to ionize and heat the IGM after matter–radiation equality \( z_{\text{eq}} \), which is much earlier than the formation of the first cosmological dark halo structure and later the first stars. We show that the dark matter annihilation density contributed by UCMHs can completely dominate over the homogeneous dark matter annihilation background, even for a tiny UCMH fraction \( f_{\text{UCMH}} = \Omega_{\text{UCMH}}(z_{\text{eq}})/\Omega_{\text{DM}} \geq 10^{-15} (1+z)^2 (m_\chi c^2/100 \text{GeV})^{-2/3} \) with a standard thermal-relic dark matter annihilation cross-section, and can provide a new gamma-ray background in the early Universe. UCMH annihilation becomes important to IGM evolution for approximately \( f_{\text{UCMH}} > 10^{-6} (m_\chi c^2/100 \text{GeV}) \). The IGM ionization fraction \( x_{\text{ion}} \) and gas temperature \( T_m \) can be increased from the recombination residual \( x_{\text{ion}} \sim 10^{-4} \) and adiabatically cooling \( T_m \propto (1+z)^2 \) in the absence of energy injection, to a maximum value of \( x_{\text{ion}} \sim 0.1 \) and \( T_m \sim 5000 \text{ K} \) at \( z \geq 10 \) for the upper bound on UCMH abundance constrained by the cosmic microwave background optical depth.

A small fraction of UCMHs are seeded by PBHs. The X-ray emission from gas accretion on to PBHs may totally dominate over dark matter annihilation, and may become the main cosmic ionization source for a PBH abundance \( f_{\text{PBH}} = \Omega_{\text{PBH}}/\Omega_{\text{DM}} \gg 10^{-11} (10^{-12}) \) with PBH mass \( M_{\text{PBH}} \sim 10^{-6} \text{ M}_\odot (10^2 \text{ M}_\odot) \). However, the constraints on gas accretion rate and X-ray absorption by baryon accumulation within UCMHs, together with accretion feedback, show that X-ray emission can only be a promising source much later than UCMH annihilation at \( z < z_m \ll 1000 \), where \( z_m \) depends on the PBH masses, their host UCMHs and the dark matter particles. Also, UCMH radiation including both annihilation and X-ray emission can significantly suppress the low-mass first baryonic structure formation. The effects of UCMH radiation on baryonic structure evolution are quite small as regards the gas temperature after virialization, but more significant in enhancing gas chemical quantities such as the ionization fraction and molecular hydrogen abundance in baryonic objects.

Key words: intergalactic medium – galaxies: structure – cosmology: theory – dark matter – early Universe.

1 INTRODUCTION

Ultracompact minihaloes (UCMHs) are primordial dark matter structures that formed by dark matter accreting on to primordial black holes (PBHs) after matter–radiation equality, \( z_{\text{eq}} \sim 3100 \), or directly by collapse on to an initial dark matter overdensity produced by a small density perturbation before \( z_{\text{eq}} \), e.g. in several Universe phase-transition epochs (Mack, Ostriker & Ricotti 2007; Ricotti & Gould 2009). If the density perturbation in the early Universe exceeds a critical value \( \delta_c = (\delta \rho / \rho)_c \sim 1/3 \), this region becomes gravitationally unstable and collapses directly to form a PBH (Hawking 1971; Carr & Hawking 1974; see Khlopov 2010 for a review, and the references therein). PBHs that form with a sufficiently high mass \( \geq 10^{16} \text{ g} \) do not evaporate but begin to grow by accreting the surrounding dark matter and form a compact dark
matter halo, which will grow by two orders of magnitude in mass during the matter-dominated era (Mack et al. 2007). These haloes are the so-called ultracompact minihaloes (UCMHs), or alternatively ‘primordially laid ultracompact minihaloes’ (PLUMs). On the other hand, a small density perturbation in the early Universe $10^{-3} < \delta < \delta_c$ will form a compact dark matter overdensity instead of a PBH. Such an overdense cloud can also seed the formation of UCMHs (Ricotti & Gould 2009; Scott & Sivertsson 2009; Josan & Green 2010). Note that the initial density perturbations from inflation were just $\delta \sim 10^{-3}$ to $10^{-2}$; it is proposed that a far more viable formation method for UCMHs is by accretion on to a dark matter overdensity, which requires a much lower perturbation threshold than do PBHs. Also, the UCMHs seeded by primordial overdensities have a different profile from those seeded by PBHs (Bertschinger 1985; Mack et al. 2007).

UCMHs have recently been proposed as a new type of non-baryonic massive compact gravitational object (MACHO: Ricotti & Gould 2009) as well as gamma-ray and neutrino sources (Scott & Sivertsson 2009). UCMHs could produce a microlensing light curve, which can be distinguished from that of a ‘point-like’ object such as a star or brown dwarf, thus becoming a promising new target for microlensing searches. Moreover, the abundance of UCMHs can be constrained by the observation of the Milky Way gamma-ray flux and the extragalactic gamma-ray background, although this constraint is still very uncertain based on today’s data (Josan & Green 2010; Lacki & Beacom 2010; Saito & Shirai 2011). Since we know the growth of an isolated UCMH as a function of redshift (Mack et al. 2007), we can natively trace the fraction of UCMHs back to very high redshift without considering mergers and tidal destruction. Until now, most work on UCMHs has focused on the properties of nearby UCMHs at $z < 1$. Another important question that has barely been discussed is the consequences of UCMH radiation at very high redshift, since UCMHs are ‘remnants’ originally from the early Universe. As sources of heating and ionization before the first structures, stars and galaxies, sufficient UCMHs might play an important role in changing the chemical and thermal history of the early Universe. Our main purpose in this paper is to investigate the impacts of UCMH emission on the intergalactic medium (IGM) in the Universe’s reionization era and the first baryon structure formation and evolution that followed.

The process of reionization of all hydrogen atoms in the IGM would have been completed at redshift $z \approx 6$ (Becker et al. 2001; Fan et al. 2002). However, much earlier ionization at $z > 6$ is implied by Wilkinson Microwave Anisotropy Probe (WMAP) observations (Dunkley et al. 2009; Komatsu et al. 2009). It is commonly suggested that possible contributions to high-redshift reionization in range approximately $6 < z < 20$ are the first baryonic objects to produce significant ultraviolet light, early (Pop III and Pop II) stars and old quasars (Barkana & Loeb 2007; Wise & Abel 2008; Meiksin 2009; Volonteri & Gnedin 2009). However, it is still unclear whether quasars and first stars were sufficiently efficient to reionize the Universe. Dark matter, on the other hand, is suggested as an exotic source of ionization and heating at high redshift due to its self-annihilation or decay. It is usually proposed that weakly interacting massive particles (WIMPs) provide a compelling solution to identify the dark matter component. The mass of the dark matter particles, $m_\chi$, and the average annihilation cross-section, $\langle \sigma v \rangle$, are the two crucial parameters that affect the ionizing and heating processes. Under the thermal-relic assumption that the cross-section $\langle \sigma v \rangle \approx 3 \times 10^{-26}$ cm$^3$ s$^{-1}$ to match the observed $\Omega_{\chi} h^2 \approx 0.110$, most previous studies showed that the effects of homogeneous dark matter background annihilation or decay on the high-redshift IGM are expected to be important only for light dark matter, $m_\chi c^2 < 1$ GeV, or sterile neutrinos (Hansen & Haiman 2004; Pierpaoli 2004; Belotsky et al. 2005; Mapelli & Ferrara 2005; Mapelli, Ferrara & Pierpaoli 2006; Zhang et al. 2006; Ripamonti, Mapelli & Ferrara 2007a,b; Chluba 2010). The annihilation flux would be enhanced only after the formation of the first dark objects at $z < 60$, as dark matter becomes more clumpy (Chuzhoy 2008; Natarajan & Schwarz 2008, 2009, 2010; Belikov & Hooper 2009, 2010). However in our case dark matter is more concentrated in UCMHs, which are significantly denser than the homogeneous dark matter background, so WIMP dark matter annihilation within UCMHs may produce powerful gamma-ray sources that dominate over the homogeneous background annihilation, even though UCMHs are very rare.

A small fraction of UCMHs are seeded by PBHs (Mack et al. 2007). In this paper we call these UCMHs ‘PBH-host UCMHs’. Since PBH abundance is still uncertain for a broad range of PBH mass (Josan & Green 2009; Carr et al. 2010) we only give a qualitative estimate that the abundance of PBH-host UCMHs should be much less than that of other UCMHs. For PBH-host UCMHs, the X-ray emission from the accreting baryonic gas that flows on to PBHs may totally dominate over dark matter annihilation within the host UCMH, since the Eddington luminosity is several orders of magnitude brighter than that of annihilation from the host UCMH and the photoionization cross-section for hydrogen or helium is much larger than the Klein–Nishina or pair-production cross-section for energetic gamma-rays. It is very difficult for a ‘naked’ PBH to reach a sufficiently high accretion rate in the IGM environment (Barrow & Silk 1979; Carr 1981; Gnedin, Ostriker & Rees 1995; Miller & Ostriker 2001; Ricotti 2007; Mack & Wesley 2008; Ricotti, Ostriker & Mack 2008), but the situation will be quite different when PBHs are surrounded by UCMHs. The accretion rate and X-ray luminosity of baryons can change significantly when the effects of a growth UCMH are involved (Ricotti 2007; Ricotti et al. 2008). However, it is possible that the gas is heated and piled up around the PBH if the host UCMH is sufficiently massive. Also, accretion feedback such as outflows or radiation pressure prevents gas from being totally eaten by the PBH immediately, if the gas accretion rate significantly exceeds the Eddington limit. As a consequence, the gas density and temperature within the UCMH may be significantly higher than the cosmic universal gas density and the X-ray emission is totally absorbed by the UCMH but reradiated basically in the infrared band. In this paper we will give criteria for X-ray emission escaping from the host UCMH to ionize the IGM. We will also compare the importance of X-ray emission from PBH-host UCMHs and dark matter annihilation from total UCMHs in the early Universe, depending on the abundance of both total UCMHs and PBH-host UCMHs.

Another topic related to UCMH radiation is that the formation and evolution history of the first baryonic structure can be changed by UCMH radiation. Previous studies showed that the annihilation or decay of extended distributed dark matter in the first structures changes both the gas temperature and the chemical properties such as the abundance of molecular coolants like H$_2$ and HD (Biermann & Kusenko 2006; Sasialak, Biermann & Kusenko 2007; Ripamonti et al. 2007b). Higher coolant abundances help to decrease the gas temperature and favour an early collapse of the baryon gas inside the halo, but dark matter energy injection delays this collapsing process. It is still under debate whether dark matter annihilation or decay inside the dark halo will promote or suppress the first structure formation. Nevertheless, it is concluded that the promotion or suppression effect is quite small for most dark matter models, as the change of gas temperature in a virialized halo for various dark
This paper is organized as follows. In Section 2, we calculate the dark matter annihilation luminosity from UCMHs and the X-ray emission from PBH gas accretion. We emphasize the importance of UCMH annihilation compared with homogeneous dark matter background annihilation, and focus on the physical reasons as to whether and when the X-ray emission from PBHs becomes more important than UCMH annihilation in the early Universe. In Section 3 we discuss the gas heating and ionization process arising from the two types of UCMH radiation from $z \sim 1000$–10 and investigate the impact of UCMH radiation on IGM evolution. Next, in Section 4 we show the influence of UCMH radiation on the first baryonic structure formation and evolution. The main results of this paper are given in Sections 3 and 4. In Section 5 we discuss the importance of UCMHs in the reionization era and the effects of a single massive UCMH on the first baryonic structure as well as other secondary effects. The reader could skip this section and go directly to Section 6, which presents the conclusions. In this paper we do not consider dark matter decay, which should have similar consequences to the annihilation process. Also, we fix the annihilation cross-section to be the thermal-relic $\langle \sigma v \rangle = 3 \times 10^{-26}$ cm$^3$ s$^{-1}$, although much larger cross-sections $\langle \sigma v \rangle = 3 \times 10^{-24}$ cm$^3$ s$^{-1}$–$10^{-20}$ cm$^3$ s$^{-1}$ are proposed in the hope of explaining the reported Galactic cosmic ray anomalies as the results of dark matter annihilation (Aharonian et al. 2008; Chang et al. 2008; Abdo et al. 2009). A larger cross-section with the same dark matter particle mass $m_\chi$ can have higher luminosity and more significant influence on ionization and heating of the early Universe. Table 1 gives the notation and definition of some quantities in this paper.

### Table 1. Notation and definition of some quantities in this paper. In this table ‘DM’, ‘anni’, ‘lum’, ‘charac.’, ‘rad’, ‘func’ are short for dark matter, annihilation, luminosity, characteristic, radiation and function. The subscript ‘acc’ is for X-ray emission because X-rays are emitted by gas accretion on to PBHs.

| Notation | Definition | Section/Eq. |
|-----------------|-----------------|---------------|
| $m_{b}(z)$ | UCMH mass at redshift $z$ | Section 2.1, equation (1) |
| $\rho_{b}(r, z)$ | UCMH (density) profile at $z$ | Section 2.1, equation (2) |
| $R_{b}(z)$ | extent radius of UCMH | Section 2.1, equation (3) |
| $m_{b}$ | DM particle mass | Section 2.1, equation (4) |
| $\langle \sigma v \rangle$ | DM average anni cross-section | Section 2.1, equation (4) |
| $L_{\text{anni}}$ | anni lum of a single UCMH | Section 2.1, equation (5) |
| $f_{\text{UCMH}}$ | $f_{\text{UCMH}}(z_{\text{eq}}) = \Omega_{\text{UCMH}}(z_{\text{eq}})/\Omega_{\text{DM}}$ | Section 2.1, equation (7) |
| $L_{\text{anni}}$ | UCMH anni lum per volume | Section 2.1, equation (10) |
| $b_{\text{bgd}}$ | homogeneous DM anni background | Section 2.1, equation (11) |
| $\epsilon_{\text{acc}}$ | X-ray rad density from PBH-host UCMHs | Section 2.2, equation (13) |
| $\Omega_{\text{PBH}}/\Omega_{\text{DM}}$ | | Section 2.2, equation (13) |
| $r_{B}$ | Bondi accretion radius | Section 2.2, equation (19) |
| $A$ | amplification factor of the IGM $T_{m}$ | Section 2.2, equation (19) |
| $f_{B}$ | dimensional accretion rate | Section 2.2, equation (23) |
| $z_{\text{com}}$ | charac. redshift for gas accretion | Section 2.2, equation (28) |
| $\chi_{\text{ion}}(z)$ | reionized baryon fraction at $z$ | Section 3.1, equation (35) |
| $\epsilon(z)$ | energy deposition rate per volume at $z$ | Section 3.1, equation (39) |
| $E_{\gamma}$ | photon energy from anni DM | Section 3.1, equation (41) |
| $E_{\text{X}}$ | charac. energy for X-ray emission | Section 3.1, equation (41) |
| $T_{m}$ | IGM gas temperature | Section 3.1, equation (43) |
| $f_{\text{H}_{2}}$ | molecular hydrogen fraction in IGM gas | Section 3.1, equation (47) |
| $L_{\text{halo}}$ | total anni lum from a dark halo | Section 4.1, equation (4) |
| $L_{\text{UCMH}}$ | anni lum from UCMHs inside a halo | Section 4.1, equation (4) |
| $L_{\text{ext}}$ | anni lum from extended DM in a halo | Section 4.1, equation (4) |
| $\epsilon_{\text{loc, iso}}$ | energy deposited by isothermal DM halo anni | Section 4.1, equation (48) |
| $\epsilon_{\text{loc,uchi}}$ | energy deposited by UCMHs anni inside a halo | Section 4.1, equation (49) |
| $\epsilon_{\text{loc, acc}}$ | energy deposited by X-rays inside a halo | Section 4.2, equation (50) |
| $z_{\text{eq}}$ | $z_{\text{eq}}$, for PBH-host UCMHs inside/outside a halo | Section 4.2, equation (52) |
| $\epsilon_{\text{rad}}$ | energy deposition rate density by $\delta$-func SED | Section 5.3, equation (52) |
| $\delta$ | factor ratio of rad of differently distributed UCMHs | Section 5.4, equation (64) |

### 2 RADIATION FROM UCMHS

In this section we discuss two types of energy emission from UCMHs in the early Universe: dark matter annihilation and X-ray emission from the accreting baryonic gas on to PBHs. As mentioned in Section 1, the second type of emission is related to a small fraction of UCMHs that host PBHs. Generally we still call the second type of energy emission ‘UCMH radiation’; this is because a PBH is always located in the centre of its host UCMH and belongs to a PBH–UCMH system.
2.1 Dark matter annihilation

The dark matter annihilation luminosity of nearby UCMHs ($z = 0$) has been calculated recently (Scott & Sivertsson 2009; Josan & Green 2010; Lacki & Beacom 2010; Saito & Shirai 2011). We assume that UCMHs stop growing at $z \approx 10$ when structure formation has progressed deeply enough to prevent dark matter from accreting further. Now we calculate the annihilation luminosity as a function of redshift in the early Universe before $z \approx 10$ and compare the result with homogeneous background annihilation. The mass of the UCMHs accreted by dark matter radial infall is given by (Mack et al. 2007; Ricotti & Gould 2009; Scott & Sivertsson 2009; Josan & Green 2010)

$$m_{\chi}(z) = \frac{\delta m}{m_{\chi}(1 + z)} \right] \frac{1 + z_{\text{eq}}}{1 + z} ,$$

where $z_{\text{eq}} \approx 3100$ is the redshift of matter–radiation equality and $\delta m$ is the mass of the initial dark matter overdensity. The density profile in a UCMH $\propto r^{-3}$ can be written as

$$\rho_{\chi}(r,z) = \frac{(3 - \alpha)m_{\chi}(z)}{4\pi R_{\chi}^3} m_{\chi}^2 r^{-3} ,$$

where the factor $(3 - \alpha)/4\pi$ in equation (2) is obtained by normalizing the total mass inside the radius of maximum halo extent $R_{\chi}$ and $R_{\chi}$ is calculated by

$$R_{\chi}(z) \approx 0.019 \rho_{\text{c}} \left( \frac{1000}{1 + z} \right) \left( \frac{m_{\chi}(z)}{M_{\odot}} \right)^{1/3} .$$

Dark matter annihilation reduces the density in the inner region of a UCMH and makes the density in this region flat. Following Ullio et al. (2002), the UCMH power-law density distribution is truncated at the maximum density

$$\rho(r_{\text{cut}}) = \rho_{\text{max}} = \frac{m_{\chi}}{(\sigma v)(t - t_1)} ,$$

where $t \approx \frac{1}{2}(1 + z)^{-3/2}(\Omega_{\text{m,0}})^{-1/2}H_0^{-1}$ is the age of the Universe at a certain redshift $z$ and $t_1 \approx 77$ kyr is the initial age at $z_{\text{eq}}$. Thus the total dark matter annihilation luminosity within the UCMH can be calculated as

$$L_{\text{ann}} = \int_0^{R_{\chi}} 2\pi r^2 n_{\chi}^2(r) (\sigma v) m_{\chi}^2 c^2 dr$$

$$= \frac{2\pi c^2}{3} \left( \frac{2\alpha}{2\alpha - 3} \right) K^{1/2} (\sigma v)^{(3-\alpha)/\alpha} (1 + z)^{9-4\alpha/\alpha}$$

$$\times \delta m(t - t_1)^{3(2\alpha)/\alpha - 3(\alpha - 3)/\alpha} ,$$

where $K = (3 - \alpha)(4.66 \times 10^{10})^{\alpha - 3}(1 + z_{\text{eq}})^{\alpha/3}(4\pi)^{-1}$ in cgs units and $n_{\chi}(r) = \rho_{\chi}(r)/m_{\chi}$ is the dark matter particle number density. The UCMH density profile can change from a steep slope $\alpha = 3$ (Mack et al. 2007) in the outer region to $\alpha = 1.5$ (Bertschinger 1985) in the inner region, if there is a PBH in the centre of the UCMH. In particular, radial infall on to a central extended overdensity shows a profile $\rho \propto r^{-9/4}$, which is more widely used as the typical density profile for most regions in UCMHs (Ricotti & Gould 2009; Scott & Sivertsson 2009; Josan & Green 2010). Taking $\alpha = 9/4$, we have

$$L_{\text{ann}} = 36.3 \Omega_{\odot} (\sigma v)_{s}^{1/3} m_{\chi}^{1/3} (1 + z) \frac{\delta m}{M_{\odot}} ,$$

where $(\sigma v)_{s} = (\sigma v)/3 \times 10^{-26}$ cm$^3$ s$^{-1}$ and $m_{\chi,100} = m_{\chi} c^2/100$ GeV.

According to equation (6), the annihilation luminosity decreases with the evolution of the Universe, basically because the annihilation flattens the inner density profile as shown in equation (4). The left panel of Fig. 1 gives UCMH annihilation luminosity for different halo profiles $\alpha$ and dark matter particle masses $m_{\chi}$. Note that the annihilation luminosity can be much brighter for lighter dark matter particles, and a shallower density profile reduces the luminosity by a factor of a few orders of magnitude higher than that with $\alpha = 3$. Thus the UCMH mass grows by up to two orders of magnitude from $z_{\text{eq}}$ to $z \approx 10$. From now on we take $f_{\text{UCMH}}$ as $f_{\text{UCMH}}(z_{\text{eq}})$, to show the initial abundance of UCMHs at matter–radiation equality, and at the current stage $f_{\text{UCMH}}$ is taken as a parameter for simplicity.

$^{1}$In Fig. 1 the annihilation luminosities following the density profile equation (2) are somewhat overestimated for a steep profile $\alpha \approx 3$, because in this case the halo mass within the truncated radius $r_{\text{cut}}$ can no longer be neglected. Thus the normalization factor of the density profile is $\propto [\ln (R_0/r_{\text{cut}})]^{-1}$, which is different from the factor $(3 - \alpha)/(4\pi)$ in equation (2). However, as we show that a steeper UCMH profile leads to a $L_{\text{ann}}$ several orders of magnitude higher than that with $\alpha = 2.25$, the conclusion that a steeper profile gives a brighter annihilation will not be changed too much even in the limiting case $\alpha = 3$. More details of the UCMH profile are discussed in Section 5.2.

Figure 1: Left: UCMH dark matter annihilation luminosity with halo profiles $\alpha = 1.5$ (black lines), 2.25 (blue lines) and 2.9 (red lines) and dark matter particle mass $m_{\chi} = 100$ GeV (solid lines), 1 GeV (dashed lines) and 100 MeV (dotted lines). We adopt $\delta m = 1 M_{\odot}$ in this figure. Right: ratio of UCMH luminosity to homogeneous dark matter background annihilation, where we take the fraction of UCMH in the total dark matter as $f_{\text{UCMH}} = 10^{-4}$ at $z = z_{\text{eq}}$ and the lines are the same as in the left panel.
The mean free path of gamma-ray photons from an UCMH with energy $E_\gamma$ is written as

$$\lambda_{UCMH} = \frac{n_\Lambda(z)\sigma(E_\gamma)}{n_\Lambda(z)\sigma} \approx 3 \times 10^3 \text{pc} \left(\frac{1000}{1+z}\right)^3,$$

where $n_\Lambda(z) = n_\Lambda(1+z)^3$ is the atomic number density at redshift $z$. The average distance between UCMHs is estimated as

$$d_{UCMH} = \left[\frac{1}{n_{UCMH}(z)} \right]^{1/3} = \left[ \frac{\langle \delta m \rangle}{f_{UCMH}\rho_{DM}(z)} \right]^{1/3} \approx 7 \text{ pc} \frac{f_{UCMH}^{-1/3} - 1000}{1+z}.$$ (9)

We have the mean free path exceed the inter-UCMH distance $\lambda_{UCMH} \gg d_{UCMH}$, except for extremely small $f_{UCMH} \ll 10^{-12}$. Therefore cosmic UCMH annihilation also gives a uniform gamma-ray background radiation as well as that produced by homogeneous dark matter. UCMH annihilation luminosity per volume is given by

$$l_{an} = \frac{L_{an}}{\delta m} f_{UCMH}(z=q)\rho_{DM}(z = 0)(1+z)^3$$
$$= 1.4 \times 10^{-26} \text{ erg cm}^{-3} \text{s}^{-1}$$
$$\times (\sigma v)^{1/2} m_{100}^{2/3} f_{UCMH}(1+z)^3.$$ (10)

On the other hand, the energy injection rate by the self-annihilation of homogeneous dark matter background per volume is

$$l_{bgd} = \langle \sigma v \rangle \frac{\rho_{DM}^2 c^2}{2m_p}$$
$$= 3.2 \times 10^{-43} \text{ erg cm}^{-3} \text{s}^{-1} (\sigma v)_9 m_{100}^{-1} (1+z)^6.$$ (11)

We compare the radiation between UCMH and normal dark matter annihilation:

$$l_{an}/l_{bgd} = 4.5 \times 10^{12} (\sigma v)_9^{-2/3} m_{100}^{2/3} \left(1 + \frac{10}{1+z}\right)^2 \frac{f_{UCMH}}{\langle \delta m \rangle} \gg 1.$$ (12)

If $f_{UCMH} \geq 2.2 \times 10^{-15} (\sigma v)_9^{1/3} m_{100}^{-2/3} (1+z)^2$, the gamma-ray background due to dark matter annihilation is dominated by UCMH annihilation. More details depending on the density profile $\chi$ and dark matter $m_p$ can be seen in the right panel of Fig. 1.

### 2.2 Gas accretion on to PBHs

The abundance of PBHs $\Omega_{PBH}$ as a fraction of total dark matter $\Omega_{DM}$ at $z < z_{eq}$ can be parametrized as $f_{PBH} = \Omega_{PBH}/\Omega_{DM}$. We ignore the PBH growth and take $f_{PBH}$ as a constant in the matter-dominated Universe for two reasons. The first reason is that, as the accretion processes had been significantly suppressed before $z \sim 10$ due to the relative motion between PBHs and baryon gas, the PBH growth time-scale, $f_{growth} \sim f_{Salp} \simeq 5 \times 10^6 \text{ yr}$, just reaches -- or is longer than -- the age of the Universe at $z \sim 10 \sim 5 \times 10^7 \text{ yr}$. The second reason is that low-mass PBHs have a lower accretion rate while high-mass PBHs are inclined to produce outflows, which further increase the accretion time-scale and make PBH growth negligible compared with its host UCMH growth. A very similar statement to keep $f_{PBH}$ constant was also proposed in Ricotti et al. (2008). Keep in mind that the PBH abundance $f_{PBH}$ is different from the UCMH initial abundance $f_{UCMH}$ at $z_{eq}$ as mentioned in Section 2.1, because a large proportion of UCMH seeds at $z_{eq}$ should be initial primordial dark matter overdensity but not PBHs. According to density primordial perturbation theory, generally we have the relation $f_{PBH} \ll f_{UCMH}$, which will be discussed in detail in Section 2.2.1.

Much work has been done to show the effects of radiation from PBH or early black hole accretion on the thermal and ionization history of the early Universe (Barrow & Silk 1979; Gendin et al. 1995; Miller & Ostriker 2001; Ricotti 2007; Ricotti et al. 2008; Ripamonti, Mapelli & Zaroubi 2008). Our goal in this section is to focus on the importance of PBH gas accretion radiation compared with the overall UCMH dark matter annihilation. The X-ray emission from accreting PBHs may lead to a very different heating and ionization history of the early Universe, compared with dark matter annihilation. The X-ray luminosity from an individual PBH with mass $M_{PBH}$ can be written as $L_{E} = 4\pi \eta G m_p M_{PBH}/\sigma v_c \simeq 3.3 \times 10^8 \eta_{1} L_{\odot}(M_{PBH}/M_{\odot})$, with $L_{\odot}$ and $\eta_1 = 0.11$ being the Eddington luminosity and average radiation efficiency of all PBHs respectively. This X-ray luminosity is much higher than the dark matter annihilation luminosity in equation (6). Thus the X-ray radiation density in the early Universe $z > 10$ can be written as

$$l_{X} = \frac{L_{E}}{M_{PBH}} \eta f_{PBH} \rho_{DM}(z = 0)(1+z)^3$$
$$\simeq 1.3 \times 10^{-26} \text{ erg cm}^{-3} s^{-1} \eta_{-1} f_{PBH}(1+z)^3.$$ (13)

Combining equations (10) and (13), the ratio between PBH accretion luminosity and UCMH dark matter annihilation luminosity is

$$\frac{l_{an}}{l_{X}} \simeq 9(\sigma v)_9^{-1/3} m_{100}^{2/3} \left(\frac{\eta_{-1} f_{PBH}}{f_{UCMH}}\right) \left(\frac{10}{1+z}\right).$$ (14)

Since the IGM heating rate due to energy injection by PBH X-ray emission or UCMH dark matter annihilation is proportional to both the energy injection rate and the IGM cross-section for all interactions suffered by the UCMH-emitted photons in the X-ray band $(E_x)$ or annihilation emitted by gamma-ray photons $(E_\gamma)$, the importance of IGM heating by X-ray emission and UCMH dark matter annihilation can be estimated through the ratio $l_{an}/l_X$. The high-energy photon interaction (or accretion) can be written as $\sigma_{tot} = \epsilon_{acc} f_{ucmh}$ (see Section 2.3 for more accurate calculations), the energy deposition in the IGM due to gas accretion $\epsilon_{acc}$ and annihilation $\epsilon_{ann}$ is estimated as

$$\frac{\epsilon_{acc}}{\epsilon_{ann}} \sim 3.2 \times 10^3 (\sigma v)_9^{-1/3} m_{100}^{2/3} \left(\frac{\eta_{-1} f_{PBH}}{f_{UCMH}}\right) \left(\frac{10}{1+z}\right).$$ (15)

which gives the first conclusion that X-ray heating may become totally dominant over dark matter annihilation in the early Universe if the PBH abundance exceeds a critical value,

$$\frac{\eta f_{PBH}}{f_{UCMH}} \geq 3.1 \times 10^{-7} (\sigma v)_9^{1/3} m_{100}^{-2/3} \left(\frac{1+z}{10}\right).$$ (16)

Theoretically the value of $\eta f_{PBH}/f_{UCMH}$ includes many uncertainties. In general, there are at least three reasons to have a low value $f_{PBH}/f_{UCMH} \ll 1$: density perturbation scenarios prefer a low initial value of $f_{Salp}/f_{UCMH}$; inefficient radiation $\eta \ll 1$ is favoured by low-mass PBHs while accretion feedback decreases $\eta$ for high-mass PBHs or PBHs with high-mass UCMHs; and X-ray emission from PBHs can be trapped inside the surrounding host UCMHs, which accumulate baryons.

### 2.2.1 PBH abundance

Either PBHs or UCMH overdensity seeds are produced by density perturbations in the very early Universe during some special epochs.
such as inflation or phase transitions. The cosmological abundance of UCMHs can be estimated by integrating from the overdensity seed threshold $-10^{-3}$ to the PBH formation threshold $\delta_c \sim 1/3$ (Ricotti & Gould 2009; Scott & Sivertsson 2009). Similarly, the PBH abundance is estimated by integrating the perturbation above $\delta_c \sim 1/3$ (Green & Liddle 1997; Green, Liddle & Riotto 1997). Assuming a Gaussian perturbation at a formation redshift $z_f \gg z_{eq}$ produces both PBHs and UCMH seeds, the ratio $f_{PBH}/f_{UCMH}$ at matter–radiation equality can be directly traced back to formation time $z_f$ (Carr et al. 2010; Khlopov 2010). As a result, the relative abundance of PBHs to UCMH overdensity seeds formed at redshift $z_f$ can be written as

$$f_{PBH} = \frac{\int_{-\infty}^{1} \exp \left( \frac{-\delta^2}{2\sigma(z_f)^2} \right) d\delta}{\int_{-\infty}^{0} \exp \left( \frac{-\delta^2}{2\sigma(z_f)^2} \right) d\delta} \simeq \exp \left( \frac{10^{-6} - \delta^2}{2\sigma^2} \right) \left( 1 - \frac{1}{18\delta^2} \right).$$

The perturbation variance at $z_f$ is roughly given by $\sigma(z_f) \simeq 9.5 \times 10^{-5}$ (Green & Liddle 1997), with $\sigma_{PBH}$ and $\sigma_{UCMH}$ for lower mass $\delta < 1$ (Lidsey, Carr & Gilbert 1995), the ratio $f_{PBH}/f_{UCMH}$ from a Gaussian perturbation is a function of horizon mass:

$$f_{PBH} \leq \exp \left( \frac{M_{h}(z_f)}{5.5 \times 10^{10} \text{g}} \right)^{(a-1/2)},$$

which means that the value of $f_{PBH}/f_{UCMH}$ becomes $\ll 1$ for $M_{h}(z_f) > 5.5 \times 10^{10} \text{g}$, not to mention the fact that the masses of dark matter overdensity seeds or PBHs are even lower than the horizon mass, $\delta m \ll \sigma_{h}(z_f)$ and $M_{PBH} \ll M_{h}(z_f)$. Combining equations (15) and (18), X-ray emission from gas accretion hardly becomes the dominant heating source in the early Universe, except for low-mass PBHs $M_{PBH} \ll M_{h}(z_f) < 3.7 \times 10^{16} \text{g}$ in the Gaussian perturbation scenario. However, PBHs in this mass range should either have disappeared within a Hubble time due to Hawking evaporation or be too small to accrete IGM gas.

As a result, the initially Gaussian density perturbation at a certain epoch is not able to generate sufficiently abundant PBHs to dominate over the total UCMH dark matter annihilation emission, basically because the large-amplitude part of the Gaussian distribution is highly suppressed. On the other hand, a non-Gaussian perturbation may give an even lower probability of PBH formation is highly suppressed. On the other hand, a non-Gaussian perturbation may give an even lower probability of PBH formation. For example, Ricotti & Gould (2009) require the host UCMHs around PBHs to have similar initial perturbation amplitude to PBHs, while Scott & Sivertsson (2009) have a less strict requirement of $\delta > 10^{-3}$ to form the initial dark matter overdensity. There are more physical uncertainties involved in estimating the abundance of PBHs and UCMHs produced by other mechanisms than a simple Gaussian distribution assumption. Therefore we still take $f_{PBH}$ as a free parameter sating

$f_{PBH} \ll f_{UCMH}$ to describe the relative abundances between PBHs and UCMHs.

### 2.2.2 Inefficient radiation

Another effect through which to constrain the X-ray luminosity density in the early Universe by gas accretion on to PBHs is the low radiation efficiency resulting from the low accretion rate on to low-mass PBHs, or the significant radiative feedback, thermal outflow and suppressed accretion rate from accretion on to high-mass PBHs or PBHs with high-mass host UCMHs.

In principle the mass distribution of PBHs is broad enough to cover the range from the Planck mass $-10^{-5}$ g to thousands of solar masses, $10^5 \text{M}_\odot$ (e.g. Carr et al. 2010). As mentioned in Section 2.2.1, if PBHs are formed from a Gaussian perturbation with variation $\sigma \propto M^{(a-1)/4}$ and index $a > 1$, low-mass PBHs should be more abundant because of the higher density perturbation variance $\sigma$ for lower mass $M$. Also, phase-transition models give a PBH mass or UCMH seed of less than $1 \text{M}_\odot$ (Scott & Sivertsson 2009). On the other hand, we should keep in mind that in an IGM environment a ‘naked’ PBH without a host UCMH can never reach the Eddington accretion rate $M_{PBH} = \dot{L}_{\text{Edd}}/c^2 \approx 1.4 \times 10^{17}(\text{PBH}/\text{M}_\odot) \text{g s}^{-1}$, unless its mass is $M_{PBH} > 360 \text{M}_\odot[1000(1+z)^3]/2$. The surrounding host UCMH increases the accretion rate if the PBH mass is $M_{PBH} > 100 \text{M}_\odot$ (Ricotti et al. 2008, their fig. 4). Note that an ideal case is $\eta \approx 0.01 m^2$ for the spherical case (Shapiro 1973a,b), which gives a much lower efficiency than the high accretion rate for $\eta \ll 0.1$.

If high-mass PBHs ($100 \text{M}_\odot < M_{PBH} < 10^5 \text{M}_\odot$) successfully form with an appreciable abundance compared with low-mass PBHs, as discussed by some previous authors (Mack et al. 2007; Saito, Yokoyama & Nagata 2008; Frampton et al. 2010), or the host UCMHs are massive than the PBHs, $M_{UCMH} > M_{PBH}$ (Ricotti & Gould 2009, for more details see Section 2.2.3), the Bondi accretion rates on to these PBHs with their host UCMHs can significantly exceed the Eddington limit after some critical redshift (Ricotti et al. 2008). However, spherical super-Eddington accretion is generally unstable and inclined to drive high mass-loss rates for thermal outflows (e.g. Smith & Owocki 2006). Recent simulations show that radiative feedback may become important to reduce or even quench the accretion process periodically (Milosavljević, Couch & Bromm 2009a; Milosavljević et al. 2009b; Park & Ricotti 2011). Also, thermal heating by the outflow energy or radiative feedback will increase the temperature of the gas around PBHs and decrease the Bondi radius and accretion rate on to PBHs. Besides the spherical accretion case, the falling gas angular momentum will become important for $m \gg 1$ and will form an accretion disc around PBHs. However, the physics of super-Eddington accretion discs is still not clearly known. Various types of super-Eddington accretion-disc models have been proposed, such as optically thick advection-dominated accretion flow (ADAF: Narayan & Yi 1994; Narayan, Mahadevan & Quataert 1998), adiabatic inflow–outflow (ADIO: Blandford & Begelman 1999), convection-dominated accretion flow (CDAF: Narayan, Igumenshchev & Abramowicz 2000), the ‘Polish doughnuts’ torus (Abramowicz, Jaroszyński & Sikora 1978) and thick slim disc (Abramowicz et al. 1988). In most cases the super-Eddington accretion disc advertcs most of its heating energy inward into the black hole without emission, and has a low radiation efficiency $\eta$. 

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for a high accretion rate $\eta < 1$ (Abramowicz et al. 1988; Narayan et al. 1998; see Abramowicz & Fragile 2011 for a review).

In a brief summary, low average radiation efficiency $\eta$ in equations (13)–(15) is favoured because of the low accretion rate on to low-mass PBHs and radiative or viscous feedback and outflows of accretion on to high-mass PBHs or PBHs with high-mass host UCMHs, which also leads to a low value of $\eta / m_{PBH}/UCMH$ and supresses the importance of PBH X-ray radiation from PBHs compared with the overall UCMH dark matter annihilation. From now on we consider that the X-ray emission is mainly contributed by accreting PBHs with $m > 1$.

2.2.3 Radiation trapping in host UCMHs

Some previous works discussed the fact that the accretion flow around PBHs is Compton-thin in most cases, since in the sub-Eddington accretion case the spherical flow is transparent near the PBH, while in the super-Eddington accretion case accretion flows are inclined to form an accretion disc (e.g. Ricotti et al. 2008). However, sufficiently high-mass UCMHs can accrete and thermalize baryons from the ambient IGM, even when there are no PBHs in the centre of these UCMHs. As the gravity potential at the outer edge of the host UCMH is mainly contributed by the UCMH mass, but the accretion on to the centre PBHs is according to the PBH mass, the accretion rate on to host UCMHs is not necessarily equal to the accretion rate on to centre PBHs. In other words, baryons can be first accumulated and virialized inside the host UCMH during accretion from the IGM to the inner UCMH region, followed by a secondary accretion on to the centre PBH and feedback (outflow) from the accreting PBHs. Based on this consideration, the baryons inside the UCMH can be divided into two components: the piled-up baryons inside the UCMH and the accretion spherical flow or disc around the centre PBH. Although the optical depth of the accretion gas or disc, which is mainly contributed by the depth around the inner horizon region $r \sim R_{\text{sch}}$, is transparent to X-ray photons, the X-ray emission can still be trapped and absorbed by piled-up baryons inside the host UCMH and can reradiate photons with much longer wavelength into the outer IGM environment. Quantitative analysis is given as follows. Part of the treatment is similar to an analogous discussion on dark matter structure formation and the baryon-filling process (Hoef1 et al. 2006; Okamoto, Gao & Theuns 2008).

If UCMH dark matter annihilation does not change the IGM temperature evolution, the IGM temperature is approximately coupled with the cosmic microwave background (CMB) temperature before the decoupling time $z_{\text{dec}} \sim 100$, and the IGM sound speed before $z_{\text{dec}}$ is $c_s \approx 5.7 \text{ km s}^{-1}[(1 + z)/(1000)]^{1/2}$. In general, UCMH annihilation heating and PBH emission without trapping increase the IGM temperature. We introduce an amplification factor $A$ such that $T_m = A T_{\text{CMB}}$ at $z > z_{\text{dec}}$, where $T_m$ and $T_{\text{CMB}}$ are the temperature of the IGM and CMB respectively and $A$ depends on the UCMH profile and annihilation properties, as we will calculate in Section 3. The sound speed $c_s \propto T^{1/2}$ becomes $c_s \approx 5.7 \text{ km s}^{-1}T_{\text{CMB}}^{1/2}[(1 + z)/(1000)]^{1/2}$ and the Bondi accretion radius (i.e., the accretion sonic sphere) of a PBH–UCMH system at $z > z_{\text{dec}}$ is

$$r_{\text{B}} \approx \frac{G m_{\text{PBH}}}{c_s^2} \approx \frac{400 \text{ pc}}{(1 + z)^2} A^{-1} \left( \frac{\delta m}{M_{\odot}} \right).$$

Equation (19) is derived under the assumption that the Bondi radius is larger than the UCMH size $r_{\text{B}} > R_{\text{B}}$. Furthermore, if $r_{\text{B}} > 2R_{\text{B}}$, the virial temperature of the host UCMH,

$$T_{\text{vir}} \approx \left( \frac{\mu m_{\text{p}}}{2 k_B} \right) \left( \frac{G m_{\text{PBH}}}{R_{\text{B}}} \right),$$

is greater than the temperature of the ambient IGM gas, $T_{\text{vir}} > T_{\text{ig}}$. According to the general virial theorem, the thermal pressure of the gas due to virialized heating is weak compared with the gravity of the UCMH. In this case we consider that IGM baryons should fall into the UCMH unimpeded, regardless of the centre PBH mass (Hoef1 et al. 2006; Okamoto et al. 2008).

The criterion $r_{\text{B}} > 2R_{\text{B}}$ at $z > z_{\text{dec}}$ gives

$$\left( \frac{\delta m}{M_{\odot}} \right) > 1600 A_{3/2}^{5/2} \left( \frac{1 + z}{1000} \right).$$

Note that higher IGM temperature around a UCMH, i.e. higher $A$, gives a higher minimum UCMH mass to attract baryons. A similar result can be derived for the case after decoupling, $z < z_{\text{dec}}$, where the IGM gas temperature decouples from the CMB temperature and drops adiabatically as $T_{\text{ig}} \propto (1 + z)^2$ without any heating sources. We still take the factor $A \geq 1$ to measure the IGM temperature increase due to annihilation, $T_m = A T_{\text{CMB}}$. Then, using the criterion $r_{\text{B}} > 2R_{\text{B}}$, we find that baryons fall into UCMHs unimpeded at $z < z_{\text{dec}}$ if

$$\left( \frac{\delta m}{M_{\odot}} \right) > 160 A_{3/2}^{5/2} \left( \frac{1 + z}{1000} \right).$$

As a result, if the UCMH initial overdensity seed is $\delta m > 1600 A_{3/2}^{5/2} M_{\odot}$ for $z_{\text{dec}} < z < 1000$, or $\delta m > 160 A_{3/2}^{5/2} M_{\odot}$ for $z < z_{\text{dec}}$, the IGM gas can always fill the UCMH no matter whether it includes a PBH or not. Otherwise, for a lower $\delta m$, the critical redshift $z_{\text{crit}}$ below which the UCMH accretes is $(1 + z_{\text{crit}}) < 0.63 A_{3/2}^{5/2} (\delta m / M_{\odot})$ for $z > z_{\text{dec}}$, and $(1 + z_{\text{crit}}) < 13 A_{3/2}^{5/2} (\delta m / M_{\odot})^{2/5}$ for $z < z_{\text{dec}}$. If the UCMH hosts a PBH in the centre, baryons are still able to pile up and become thermalized in the host UCMH due to gas virialization.

The lower bound of the gas accretion rate into the UCMH can be estimated as

$$M_{\text{UCMH}} = 4\pi G^2 c_s^4 \rho_{\text{igm}}(z) > 4 \pi R_{\text{vir}}^2 v_T(R_{\text{B}}) m_{\text{igm}}(z),$$

$$\sim 8.2 \times 10^{16} \text{ g cm}^{-3} (1 + z)^{1/2} \left( \frac{\delta m}{M_{\odot}} \right),$$

where $v_T(R_{\text{B}})$ is the free-fall velocity at $R_{\text{B}}$. If all the gas in the UCMH is totally accreted on to the centre PBH, the dimensionless accretion rate of the PBH $m = M_{\text{ucm}} / M_{\text{B}}$ is

$$m > 18.0 \left( \frac{1 + z}{1000} \right)^{1/2} \left( \frac{\delta m}{M_{\odot}} \right) > 1,$$
CDAFs may produce a ‘convective envelope’ with no accretion on to the black hole (Narayan et al. 2000). In general, accretion discs with super-Eddington accretion rate are inevitably accompanied by outflows and winds, which significantly decrease the final accretion rate on to the black hole. In the PBH case, these outflows should be injected back to the host UCMH environment.

The upper bound of the baryonic fraction in the UCMH is the universal fraction $\Omega_b/\Omega_m$. However, as the UCMH grows following equation (1), we adopt a more conservative method to estimate the lower bound of baryon fraction $f_b$ inside the UCMH. We estimate the baryonic fraction in a UCMH, $f_b$ (the mass ratio between gas and dark matter), as

$$ (M_{\text{UCMH}} - M_{\text{PBH}})(t - t_i) \sim m_i(z)f_b, $$

where we take $M_{\text{UCMH}} \gg M_{\text{PBH}}$, i.e. most of the gas accreted into the UCMH is piled up without being immediately eaten by the PBH. Combining equations (22) and (24), we have a lower bound on the baryon fraction of $f_b \geq 7.6 \times 10^{-3}$, which is a constant independent of the redshift.

The optical depth of the piled-up gas in the UCMH due to Compton scattering is

$$ \tau \sim x_c\sigma_T \int \frac{\rho_g(r) f_b}{\mu m_p} \, dr. $$

Since $f_b$ from equation (24) is a constant, and UCMH growth does not change the steep region $\rho_{PM} \propto r^{-\alpha}$ with $\alpha = 9/4$ but only increases $R_b$ and flattens the region $r < R_{\text{out}}$ (see equation [4]), we take the baryon fraction to be uniformly distributed in the UCMH, in both steep and flat regions. Therefore the column density of the baryon gas inside the UCMH depends on the UCMH profile, which depends on the dark matter properties ($\sigma v$, $m_i$) given by equation (4). Actually the baryon profile can be steeper in the flat region of the halo, $r < R_{\text{cut}}$, since the dark matter annihilation flattens the inner halo profile, thus giving an even larger optical depth. Furthermore, we consider that the baryonic gas is ionized, $x_e \sim 1$, inside the UCMH, at least in the flat region $r < R_{\text{cut}}$. We check that the Strömgren radius of the PBH emission $r_S$ satisfies $r_{\text{out}} < r_S < R_b$, as the heated gas near the PBH can reach a temperature as high as the Compton temperature, $\sim 10^4 \text{eV}$ in the ionized region. The hot ionized gas around the PBH produces a small sonic sphere in the dense baryon region near the PBH because the accretion rate on to the PBH, giving $M_{\text{UCMH}} \gg M_{\text{PBH}}$ as mentioned in equation (24). Note that there should be two distinct sonic spheres, the sphere for the host UCMH outside $R_b$, and that for the central PBH inside the UCMH. This scenario is similar to that of Wang, Chen & Hu (2006) in which an accreting BH has two Bondi spheres, a smaller inner sphere in the hot gas region and a larger one in the outer cooler region. Therefore we take $x_e \sim 1$ in equation (25). The optical depth is written as

$$ \tau \sim \left( \frac{\sigma v f_b}{\mu m_p} \right) \int \rho_g(r) \, dr \geq 10 \left( \frac{\delta m}{M_\odot} \right)^{1/3} \left( \frac{1 + z}{1000} \right)^{5/6} m_{\chi,100}^{9/4}(\sigma v)^{-13/9}. $$

Hereafter we take $(\sigma v)_\chi = 1$. Combining equations (20) and (26), we conclude that before the decoupling $z > z_{\text{dec}}$ the gas is always Compton-thick to X-ray emission from the centre PBH accretion when the host UCMH itself accretes baryons. After decoupling, $z < z_{\text{dec}}$, the redshift range in which X-ray emission escapes is

$$ \left\{ \begin{array}{ll}
1 + z < 13 \left( \frac{\delta m}{M_\odot} \right)^{2/5} \left( \frac{\delta m}{M_\odot} \right)^{3/5} m_{\chi,100}^{2/3} & < 7 m_{\chi,100}^{-5/2} A^{3/4}, \\
1 + z < 62 \left( \frac{\delta m}{M_\odot} \right)^{2/5} m_{\chi,100}^{2/3} & > 7 m_{\chi,100}^{-5/2} A^{3/4}.
\end{array} \right. $$

From equation (27) there is a maximum $z_{\text{cut}}$ in the case of $r_b > 2 R_b$:

$$ z_{\text{cut}}^{(r_b < 2 R_b)} = \frac{29 m_{\chi,100}^{-1/3} A^{3/4}}{10}. $$

If $z > z_{\text{cut}}$, X-ray photons will be totally trapped. Note that $z_{\text{cut}}$ sensitively decreases with the increase of IGM temperature factor $A$. A larger optical depth due to a steeper baryon profile at $r < R_{\text{cut}}$ gives an even lower $z_{\text{cut}}$. Also, the range of $\delta m$ applied in equation (27) is $0.64 m_\odot A^{3/4} < \delta m < 80 m_\odot m_{\chi,100}^{-5/3}$.

In other words, in the $r_b > 2 R_b$ case no X-rays can escape the host UCMH if $\delta m \geq 80 m_\odot m_{\chi,100}^{-5/3}$. More massive PBHs find it easier to reach the Eddington accretion rate, but more difficult to produce a transparent baryon environment in the host UCMHs.

On the other hand, if $r_b < 2 R_b$ (i.e. $T_{\text{in}}(R_b) < T_{\text{in}}$), most of the UCMH gravity potential well is not deep enough to compress the gas and overcome the pressure barrier of gas virialization heating. In this case the UCMH itself cannot accrete and thermalize baryons except for the region within radius $r_b^*$, which satisfies $[G m_b(r < r_b^*)] = c^2$. At $z > z_{\text{dec}}$, the critical radius $r_b^*$ for a pure UCMH profile is $r_b^* \propto r^{-\alpha}$.

$$ r_b^* = R_b \left( \frac{G m_b}{2 c^2} \right)^{1/(\alpha - 2)}, $$

where we apply $\alpha = 9/4$ from equation (5). The part of the UCMH inside $r_b^*$ can accrete and heat baryons. Similarly to equations (22) to (26), the accretion rate in region $r \leq r_b^*$ is

$$ \dot{m} = \frac{2 \pi c^2 \mu m_b n_b(r,z)}{M_{\text{PBH}}} \approx 3.9 \times 10^{-11} A^{-15/2} \left( \frac{100}{1 + z} \right)^{9/2} \left( \frac{m_{\chi,100}}{M_\odot} \right)^{5/4} \left( \frac{\delta m}{M_\odot} \right)^{1/2}. $$

where the factor $2 \pi$ is due to the suppressed accretion at $r > r_b^*$ in the UCMH; thus the baryon density is half of the ambient gas density. Assuming $\delta m > M_{\text{PBH}}$, equation (31) shows that only high-mass PBHs ($M_{\text{PBH}} > 100 M_\odot$) are able to produce super-Eddington accretion if there is no accretion feedback. This is basically consistent with the results in Ricotti et al. (2008). However, we should mention two points. The first point is that, if $\delta m > M_{\text{PBH}}$, the ideal accretion rate in equation (31) also increases. The second point, which is similar to the analysis below equation (23), is that the real accretion rate on to the PBH is lower than the ideal $\dot{m}$ due to higher gas temperature and accretion feedback. As a result, we find that baryons can be accumulated and virialized inside the UCMH region $r \leq r_b^*$ with a baryon fraction of approximately $f_b \propto 10^{-3} A^{-9/2}(1 + z)^{-3/4} (\delta m / M_\odot)^{1/2}$. Using the condition $\dot{m} > 1$ in equation (31) at $z > z_{\text{dec}}$ for sufficient radiation efficiency, the baryon optical depth inside radius $r_b^*$ is

$$ \tau \sim 1.1 \times 10^{-2} A^{-9/2}(1 + z)^{-13/6} m_{\chi,100}^{5/6} \left( \frac{\delta m}{M_\odot} \right)^{10/3} > 5.0 m_{\chi,100}^{5/6} A^{1/2} \left( \frac{M_{\text{PBH}}}{\delta m} \right)^{2/3} \left( \frac{1 + z}{100} \right)^{5/6} > 1. $$

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where we also take the baryon profile inside the UCMH as proportional to the dark matter profile for simplicity, and $c_{\text{H}} \sim 1$. According to equation (32), we consider that the UCMH is optically thick to X-ray emission in the case of $r_{\text{BH}} < 2R_\text{BH}$ and $z > z_{\text{dec}}$, unless $\delta m \gg M_{\text{BH}}$ or dark matter particle mass $m_{\chi} \gg 1$.

After decoupling, $z < z_{\text{dec}}$, we find that baryons can be accumulated and virialized inside a radius

$$ r_{\text{BH}} \approx 0.012 \, \text{pc} \, A^{-4} \left( \frac{30}{1 + z} \right)^6 \left( \frac{\delta m}{M_\odot} \right)^3, $$

with upper bound $z_{\text{m}}$, such that

$$ z_{\text{m}} \left( \frac{R_{\text{BH}}}{r_{\text{BH}}} \right) \approx 32 A^{-3/8} \left( \frac{\delta m}{M_{\text{BH}}} \right)^{1/2} \left( \frac{m_{\text{BH}}}{100} \right)^{-5/12} - 1, $$

below which ($z < z_{\text{m}}$) the gas is Compton-thin inside radius $r_{\text{BH}}$.

A higher ratio $\delta m/M_{\text{BH}} \gg 1$ or lighter dark matter $m_{\chi} \gg M_{\text{BH}}$ increases $z_{\text{m}}$, which also decreases slightly if the annihilation effect is included to heat the IGM gas ($A > 1$).

Not only can baryons inside UCMHs trap X-ray photons, but also outflows driven by PBH accretion feedback absorb X-ray emission. As mentioned before, the super-Eddington accretion rate on to the PBH is unstable to trigger strong outflows in both spherical and disc cases. An optically thick ‘outflow envelope’, in both polar and equatorial regions around the PBHs, forms to cover the PBH and totally or mostly absorbs X-ray emission from the inner accretion flows (Igumenshchev, Narayan & Abramowicz 2003; Kohri, Narayan & Piran 2005; Poutanen et al. 2007; Abolmasov, Karpov & Kotani 2009). In this case, the Compton heating in the outflow region should be important, in order to increase the gas pressure and temperature, balance the gravity well, reemit thermalized photons from outflows and regulate the accretion rate on to the PBH (Wang et al. 2006). It is likely to have a steady-state or periodically changing outflow envelope covering the whole PBH, but the details are still an open question beyond the purpose of this paper. What we want to show is that, even though the accretion disc itself around a PBH is optically thin to X-ray radiation, X-ray emission can still be absorbed in the outflow envelope due to accretion feedback and disc instability in the super-Eddington case, not to mention optically or geometrically thick discs, which absorb X-ray emission by themselves.

We give a summary of Section 2.2. Equation (13) computes the X-ray radiation density due to baryon gas accretion on to PBHs. The ratio $\eta_{\text{PBH}}/\eta_{\text{UCMH}}$ is parametrized in this paper. The much lower probability of PBH formation compared with UCMH formation and inefficient radiation due to low-mass PBHs ($\delta m \ll 100 M_\odot$) or accretion feedback from high-mass PBHs or PBHs with massive UCMHs ($\delta m \gg M_{\text{PBH}}$) give $\eta_{\text{PBH}}/\eta_{\text{UCMH}} \ll 1$. Moreover, we simply introduce a critical redshift $z_{\text{dec}}$ below which ($z < z_{\text{dec}}$) X-ray emission from super-Eddington accretion PBHs ($m_{\chi} \gg 1$) becomes important to heat and ionize the early Universe. A hotter IGM heated by other energy sources (e.g. annihilation) slightly decreases $z_{\text{dec}}$. Ricotti et al. (2008) showed that the accretion becomes $m_{\chi} \gg 1$ at $z_{\text{dec}} \sim 20 (100)$ for $M_{\text{PBH}} = 10^2 (300) M_\odot$, and $m_{\chi}$ is always $m_{\chi} \gg 1$ ($m_{\chi} \gg 1$) for $M_{\text{PBH}} > 10^5 M_\odot (M_{\text{PBH}} \gg 10^5 M_\odot)$. However, as we discussed in Section 2.2.3, the stage of super-Eddington accretion heating the IGM can be delayed, because X-ray photons are trapped inside the total or inner region of the UCMHs due to the accretion and virialization of the accretion gas. Outflows can also (partly) absorb X-rays. We take the critical redshift $z_{\text{dec}}$ as shown in equations (28) and (34); for $z < z_{\text{dec}}$, X-ray photons from super-Eddington accretion PBHs could escape their host UCMHs. If the PBH abundance is much less than that of UCMHs, but still satisfies equation (16), X-ray emission should dominate over early dark matter annihilation at $z < z_{\text{dec}}$.

### 3 REIONIZATION AND HEATING OF THE IGM

#### 3.1 Basic equations

The evolution of baryon ionization fraction $x_{\text{ion}}$ is given by the differential equation (e.g. Cirelli, Iocco & Panci 2009)

$$ \frac{dx_{\text{ion}}(z)}{dt} = I(z) - R(z), $$

where $I(z)$ and $R(z)$ are the ionization and recombination rates per volume respectively and $n_A$ is the atomic number density today. We use the rate $R(z)$ from Natarajan & Schwarz (2008). The ionization rate per volume due to dark matter annihilation or X-ray emission from gas accretion is given by

$$ I(z) = \int_{E_{\text{eq}}}^{E_z} dE_\gamma \, \frac{dn(z)}{dE_\gamma} \, P(E_\gamma, z) N_{\text{ion}}(E_\gamma), $$

where $E_\gamma = m_e c^2$ is the maximum energy of the emitted photon, $E_{\text{eq}} \approx \gamma_{\text{ion}}(1 + z)(1 + z_{\text{ion}})$ and the differential term $dn(z)/dE_\gamma$ is the photon spectral number density at redshift $z$. We follow Cirelli et al. (2009, see also Belikov & Hooper 2009; Natarajan & Schwarz 2008, 2009) to calculate the probability of primary ionizations per second $P(E_\gamma, z)$, and the number of final ionizations generated by a single photon of energy $E_\gamma$, $N_{\text{ion}}(E_\gamma)$. Note that $N_{\text{ion}}(E_\gamma)$ is proportional to the ionization factor $x_{\text{ion}}(z) \approx (1 - x_{\text{ion}})/3$ (Shull & van Steenberg 1985; Chen & Kamionkowski 2004), which means that approximately one-third of emitted energy goes into the reionization of atoms if $x_{\text{ion}} \ll 1$.

First of all, we consider the ionization is due to UCMH dark matter annihilation. The spectral number density is obtained as

$$ \frac{dn_{\text{ann}}(z)}{dE_\gamma} = \int_{\gamma_{\text{ion}}}^{\gamma_{\text{dec}}} d\gamma \, \frac{c}{dE_\gamma} \, \frac{dn(z)}{dE_\gamma} \left( \frac{1 + \gamma}{1 + z} \right)^3 \exp(-\tau), $$

where $\gamma_{\text{ann}}$ is the ionization luminosity as mentioned in Section 2.1. $E_{\gamma_{\text{dec}}}(z) = E_{\gamma}(1 + z)/(1 + z_{\text{dec}})$. The optical depth $\tau$ is

$$ \tau = \int_{\gamma_{\text{ion}}}^{\gamma_{\text{dec}}} d\gamma \, \frac{c}{dE_\gamma} \, n_A(1 + z_{\text{dec}})^3 \sigma_{\text{ion}}(E_\gamma), $$

with $E_{\gamma_{\text{ion}}} = E_{\gamma}(1 + z)/(1 + z_{\text{ion}})$. The total cross-section $\sigma_{\text{ion}}(z)$ for the DM annihilation photon to interact with electrons in the IGM mainly includes the Klein–Nishina cross-section for Compton scattering (Rybicki & Lightman 2004) and the photonization cross-section for H and He, $\sigma_{\text{H,He}}$ (Zdziarski & Svensson 1989). Pair production on matter becomes important for $m_e > 1$ GeV. CMB photons also contribute to the total cross-section for $m_e > 10$ TeV, which can be neglected in our case.

The total energy deposition per second per volume at redshift $z$ is given by

$$ \epsilon(z) = \int_{E_{\text{eq}}}^{E_z} dE_\gamma \frac{dn_{\text{ann}}(z)}{dE_\gamma} \sigma_{\text{ann}}(E_\gamma) E_\gamma. $$

If we take the monochromatic dark matter annihilation emission for simplicity, i.e. photons produced by dark matter annihilation have the rest energy of the dark matter particle $m_e c^2$, the photon flux spectral density then can be calculated using the $\delta$-function

$$ \frac{dn(z)}{dE_\gamma} \approx I_{\text{ann}}(z') \delta \left( E_{\gamma'} - E_{\gamma} \right). $$

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Thus we have an energy deposition $\epsilon(z)$ of

$$\epsilon(z) = \frac{\int_{z_0}^{z_f} \frac{dE}{E} n_A (1 + z)^3 \frac{dt}{dz_0} f_{\text{ann}}(z_0') \left( \frac{1 + z}{1 + z_0'} \right)^3 e^{-\sigma_{\text{tot}}(E_f)} \, dz_0',$$

(41)

where $z_0'$ satisfies

$$z_0' = \frac{E}{E_f} (1 + z) - 1,$$

(42)

and $r$ is calculated from $z_0'$ to $z$. Keep in mind that in the above formula (41) the dark matter annihilation products are simplified as gamma-ray photons with sole energy $E_f = m_f c^2$. A more realistic annihilation spectrum is model-dependent. For example, $dn/dE_f$ can be chosen following the model in Bergström, Ullio & Buckley (1998) and Feng, Matech & Wilczek (2001).

For X-ray photons from an accreting PBH, equation (41) will still be available if we choose $\sigma_{\text{tot}}$ as the X-ray total cross-section, $\sigma_{\text{tot}} \approx \sigma_{\text{H},\text{H}} + \sigma_r$, and change $E_f$ to be the X-ray characteristic energy $E_X$ as $E_X \approx 3 \text{ keV} (M_{\text{PBH}}/M_\odot)^{-1/4}$ (Salvaterra, Haardt & Ferrara 2005). As the real accreting PBH spectral energy distribution is very model-dependent (Shakura & Sunyaev 1973; Sazonov, Ostriker & Sunyaev 2004; Salvaterra et al. 2005; Ripamonti et al. 2008), in the very first calculation we simplify the X-ray emission as single-frequency emission at a characteristic energy $E_X$, which can mostly be considered as the peaked energy in the real spectral energy distribution. We will also discuss a more realistic PBH spectral energy distribution in Section 5.3.

Now we list the heating and cooling processes in the IGM. The heating of the IGM by UCMH annihilation or X-ray emission can be written as

$$\frac{dT_m}{dt}_{\text{ann}} = \frac{2}{3m_H n_A (1 + z)^3 (1 + f_{\text{He}} + x_{\text{ion}})} \epsilon(z),$$

(43)

where the heating fraction $\eta_{\text{heat}}(x_{\text{ion}})$, which shows the portion of energy $\epsilon(z)$ that goes into heating the IGM, is adopted as $\eta_{\text{heat}} = C |1 - (1 - x_{\text{ion}})^{a} |^b$, with $C = 0.9971$, $a = 0.2663$ and $b = 1.3163$ (Shull & van Steenberg 1985). We can approximate the He fraction in the IGM as $f_{\text{He}} \approx 0.7$. Moreover, CMB photons can be treated as another heating source for the IGM if the IGM gas is colder than the CMB ($T_m < T_{\text{CMB}}$), otherwise the IGM gas would transfer energy into the CMB environment. The coupling between IGM gas and CMB photons can be important when the difference between $T_m$ and $T_{\text{CMB}}$ is significant (Weymann 1965; Tegmark et al. 1997; Seager, Sasslov & Scott 2000):

$$\frac{dT_m}{dt}_{\text{comp}} = k_{\text{comp}} T_{\text{CMB}}^3 v_{\text{ion}}(T_{\text{CMB}} - T_m),$$

(44)

where the coupling rate constant $k_{\text{comp}} \approx 5.0 \times 10^{-22} \text{ s}^{-1}$.

Other IGM cooling terms are dominated by the adiabatic cooling during the expansion of the Universe as

$$\frac{dT_m}{dz}_{\text{ad}} = \frac{2T_m}{1 + z},$$

(45)

for low temperature. Note that the IGM temperature will decrease independently as $T_m \propto (1 + z)^3$ for a pure adiabatic cooling process. Furthermore, for sufficiently high temperature $\sim 10^4 \text{ K}$, the molecular hydrogen $H_2$ cooling will also be important. This cooling term can be calculated as

$$\frac{dT_m}{dt}_{\text{H}_2} = \Lambda_{\text{H}_2}(1 - x_{\text{ion}} - 2f_{\text{H}_2}) f_{\text{H}_2} n_A (1 + z)^3,$$

(46)

where we adopt the specific cooling coefficient $\Lambda_{\text{H}_2}$ from Hollenbach & McKee (1979) and Yoshida et al. (2006). We neglect other

c hemical cooling processes such as bremsstrahlung, helium-line cooling, $H_2$ line cooling and hydrogen three-body reaction. HD cooling is important for $T \lesssim 200 \text{ K}$ and low densities (Yoshida et al. 2006), but the gas adiabatic cooling will be dominant in this case. As we mainly pay attention to the evolution of the ionized fraction $x_{\text{ion}}$ and IGM temperature $T_m$, we only include the evolution of hydrogen $(H, H^+, H_2)$ and electron gas $(e^-)$ as the main species for ionization. More detailed simulation including other species and cooling processes is beyond the purpose of this paper. The evolution of the $H_2$ fraction is adopted from the semi-analytic model in Tegmark et al. (1997):

$$\frac{dH_2}{dt} = k_{\text{H}_2} n_A (1 + z)^3 (1 - x_{\text{ion}} - 2f_{\text{H}_2}),$$

(47)

where we follow Tegmark et al. (1997) and Galli & Palla (1998) to calculate the reaction coefficient $k_{\text{H}_2}$.

### 3.2 Solutions of IGM evolution

Since no direct evidence related to the UCMH radiation has been confirmed until now, the UCMH abundance is still uncertain. In this paper we take the UCMH fraction $f_{\text{UCMH}}$ as a free parameter. In Fig. 2 we show the ionization fraction $x_{\text{ion}}(z)$ and the IGM temperature $T_m$ for $f_{\text{UCMH}} = 10^{-4}$ and $10^{-6}$, which corresponds to today’s expected abundance $\sim 1$ per cent and $10^{-4}$ respectively. Moreover, general discussion on the UCMH abundance will be given later. The initial $x_{\text{ion}}$ at $z = 1000$ is taken as 0.01, and $T_m$ as the CMB temperature (Galli & Palla 1998; Ripamonti 2007; Ripamonti et al. 2007a,b). The basic conclusion, which is similar to the previous work, is that lighter dark matter particles or higher UCMH abundance give larger $x_{\text{ion}}$ and higher $T_m$. An extremely bright UCMH annihilation background (e.g. $m_1 c^2 \lesssim 1 \text{ GeV}$ for $f_{\text{UCMH}} \sim 10^{-4}$ or $m_1 c^2 \lesssim 100 \text{ MeV}$ for $f_{\text{UCMH}} \sim 10^{-6}$) even gives a monotonic increase in $x_{\text{ion}}$ and $T_m \gtrsim 10^4 \text{ K}$ without a standard reionization epoch in the early Universe. We compare the UCMH annihilation results with the homogeneous background annihilation; note that the homogeneous dark matter annihilation background only makes noticeable effects for light dark matter particles $m_1 c^2 < 1 \text{ GeV}$ or sterile neutrinos. UCMHs, which provide a new dominant dark matter annihilation gamma-ray background as shown in Section 2.1, play a more important role in ionizing and heating the early Universe.

Furthermore, as the cosmological Jeans mass $m_J$ can be taken as an indicator of IGM structure evolution, the left panel of Fig. 3 gives the evolution of Jeans mass in the cosmic UCMH annihilation background. The Jeans mass $m_J \propto T_m^{1/3} n^{-1/2}$ should be a constant if the gas temperature is always equal to the CMB temperature $T_m = T_{\text{CMB}}$. In this paper we call this constant ‘CMB mass’. In the left panel of Fig. 3, $m_J$ with various $T_m$ is generally normalized in units of ‘CMB mass’. The remaining Thomson scattering optical depth contributed by UCMH annihilation is shown in the right panel of Fig. 3. For $6 \lesssim z < 30$, the remaining CMB optical depth is estimated as $\delta \alpha \approx 0.046 \pm 0.016$ by the WMAP five-year measurement. We assume a linear increase of $x_{\text{ion}}$ from $z = 10$ to the full ionization

$$\text{WMAP five-year measurements give the CMB Thomson scattering optical depth $\delta \alpha \approx 0.046 \pm 0.016$ (Komatsu et al. 2009), which is mostly due to ionization at late times $z < 30$ (Ricotti et al. 2008; Natarajan & Schwarz 2010). If we subtract the optical depth contributed by the totally ionized gas $\alpha (z \leq 6) = 0.038$, the remaining depth is $\delta \alpha \approx 0.046 \pm 0.016 \pm 0.062$ (Cirelli et al. 2009).}$$
Figure 2. Effects of UCMH dark matter annihilation on IGM evolution. The upper two panels correspond to the case of $f_{\text{UCMH}} = 10^{-4}$ and the lower to $f_{\text{UCMH}} = 10^{-6}$. The thick (black) lines show, from top to bottom, the results for UCMH annihilation with $m_\chi c^2 = 100 \text{ MeV}$, $1 \text{ GeV}$, $10 \text{ GeV}$, $100 \text{ GeV}$ and $500 \text{ GeV}$, while the thin (blue) lines give the results of homogeneous dark matter background annihilation for $m_\chi c^2 = 100 \text{ MeV}$, $1 \text{ GeV}$ and $10 \text{ GeV}$.

Figure 3. Left: ratio of Jeans mass to ‘CMB mass’, which is the Jeans mass for a gas temperature always equal to the CMB temperature. Right: evolution of the remaining Thomson scattering optical depth $\delta \tau$, with the grey belt showing the WMAP 5-year 1σ of the remaining CMB optical depth at $10 \leq z < 30$. The thick (black) lines are for $f_{\text{UCMH}} = 10^{-4}$ and the thin (blue) lines for $f_{\text{UCMH}} = 10^{-6}$. The lines of the same colour, from the top down, are for $m_\chi c^2 = 100 \text{ MeV–500 GeV}$ as in Fig. 2.

time $z = 6$; thus the upper bound contribution of UCMH annihilation to the measurable CMB optical depth is $\delta \tau \simeq 0.028 \pm 0.016 \leq 0.044$. Based on this consideration, in our examples only one extreme case, $m_\chi c^2 = 100 \text{ MeV}$ with $f_{\text{UCMH}} = 10^{-4}$, is ruled out by the CMB remaining optical depth in Fig. 3. UCMH annihilation can significantly increase the Thomson optical depth in the early Universe, $z \gg 100$, up to $\delta \tau \sim 0.5$ without stringent constraints from the CMB optical depth measurement at $z < 30$. After the last scattering epoch, the ionization fraction $x_{\text{ion}}$ can change from $x_{\text{ion}} \sim 10^{-4}$ without dark matter annihilation to an upper bound $x_{\text{ion}} \sim 0.1$ (e.g. $m_\chi c^2 = 1 \text{ GeV}$ and $f_{\text{UCMH}} = 10^{-6}$). Also, the IGM can be heated from a temperature $T_m$ of adiabatically cooling $T_m \propto (1 + z)^2$ in the absence of a heating source to the upper bound $T_m > 10^3 \text{ K}$ with a sufficient amount of heating contributed by UCMH annihilation. A Jeans mass much higher than the ‘CMB mass’, an increase of $\sim 2–3$ orders of magnitude, can be obtained due to the hotter IGM temperature. Therefore, we can natively estimate that the formation of small baryonic objects can be strongly suppressed, although more investigations need to be carried out in Section 4.

In general, we find that the impact of UCMH annihilation on IGM evolution can be empirically estimated by the factor $m_\chi^{-1} \chi$, $100 f_{\text{UCMH}}$, while the threshold of UCMH abundance affecting IGM evolution is approximately given by $m_\chi^{-1} \chi, f_{\text{UCMH}} > 10^{-2}$, with an upper bound constrained by the CMB optical depth at late times $z < 30$ as $x_{\text{ion}} \sim 0.1$ and $T_m \sim 5000 \text{ K}$ at $m_\chi^{-1} \chi, f_{\text{UCMH}} \sim 10^{-2}$. The CMB optical depth enhancement at early times $z > 30$ can be more dramatic than that at late times due to the higher annihilation luminosity at early redshifts (equation 6). This is different from the PBH radiation,
which has higher luminosity at late times (Section 2.2.3; Ricotti et al. 2008). Further phenomenological constraints should be made by CMB polarization anisotropies; this is left for a future investigation. Keep in mind that another channel to concentrate dark matter rather than primordial density perturbations is the formation of the first dark objects, which should affect the IGM evolution much later ($z < 100$) than UCMHs. Therefore UCMH annihilation definitely has a much earlier and more important impact on IGM evolution from last scattering to structure formation time.

So far we have given results only for UCMH dark matter annihilation. Whether the X-ray emission from the PBH-host UCMHs will significantly change the above results depends mainly on the fraction of PBHs $f_{\text{PBH}}$, the average inflow radiation efficiency $\eta$ and the critical redshift $z_m$, as given in Section 2.2. In our paper we combine the product $\eta f_{\text{PBH}}$ as one parameter. Remember the results in Section 2.2: when X-rays from the centre PBH region successfully passes through the transparent baryon medium in the host UCMH at $z \lesssim z_m$, the much brighter X-ray luminosity and much larger interacting cross-section $\sigma_{\text{int}}(E_X)$ compared with the annihilation luminosity and $\sigma_{\text{ann}}(E_\gamma)$ usually guarantee that X-ray emission dominates over UCMH annihilation (equations 15), except for a much lower PBH fraction $f_{\text{PBH}}$ below the value in equation (16). We will give a lower limit of $f_{\text{PBH}}$ above which X-rays have obvious impact on the IGM evolution at $z \lesssim z_m$. In the following calculation we assume that equation (16) is always satisfied and do not distinguish between the redshift dividing the UCMH radiation into annihilation-dominated or X-ray-dominated and the redshift giving a transparent baryon environment in UCMHs, but simply use a single parameter $z_m$.

In Fig. 4 we take $z_m$ and $\eta f_{\text{PBH}}$ as parameters with characteristic emission frequency $E_X = 1, 10$ and $100$ keV, which corresponds to typical PBH masses of $10^2 M_\odot$, $10^{-2} M_\odot$ and $10^{-6} M_\odot$ respectively. Higher $z_m$ means a higher ratio $\delta_m/M_{\text{PBH}}$ or lighter dark matter particles. Only X-ray emission as an energy source is calculated in this figure. As shown in Fig. 4, the final properties of the IGM at $z \sim 10$ with the same $\eta f_{\text{PBH}}$ and $E_X$ are more or less close to each other regardless of the value of $z_m$, which means that lower $z_m$ gives more dramatic thermal and chemical change at $z <$...
According to equation (34), lower \( z_{\text{min}} \) corresponds to lighter \( m_{\chi} \), which also increases the UCMH annihilation. On the other hand, \( x_{\text{min}} \) and \( T_{\text{min}} \) vary by more than three orders of magnitude from \( E_X \sim 100 \text{ keV} \) (10\(^{-6}\) \( M_\odot \)) to \( E_X \sim 1 \text{ keV} \) (10\(^2\) \( M_\odot \)) with the same \( \eta_{\text{PBH}} \), which means that massive PBHs favour PBHs. Also, the X-ray radiation effect can be neglected when \( \eta_{\text{PBH}} \leq 10^{-11} \) for \( E_X \sim 100 \text{ keV} \), but a smaller limit \( \eta_{\text{PBH}} \leq 10^{-12} \) is applied for \( E_X \sim 1 \text{ keV} \). Below the lower limit, the PBHs are not expected to have any promising effects on reionization. The estimate in Fig. 5 shows that no strict constraints are made for \( \eta_{\text{UCMH}} \leq 10^{-6} \) by the remaining CMB depth \( \delta_T \), which allows a dramatically increased Jeans mass due to the hot IGM gas \( \sim 10^4 \text{ K} \).

As a result, X-ray emission from PBHs gives a more promising impact on the IGM evolution if \( \eta_{\text{PBH}} \gg 10^{-11} \) (10\(^{-12}\)) for \( M_{\text{PBH}} \sim 10^{-6} \odot \) (10\(^7\) \( M_\odot \)), or empirically \( \eta_{\text{PBH}} \gg 1.8 \times 10^{-12} (M/M_\odot)^{1/8} \). As we assume \( f_{\text{UCMH}} \leq 10^{-4} \), we expect that the UCMH dark matter annihilation only played its role on IGM evolution at very high redshift \( z_{\text{min}} \leq z < 1000 \), but X-ray emission changes the Universe reionization history dramatically at relatively lower redshift \( z < z_{\text{min}} \). Considering \( \eta \sim 0.1 \), the upper-bound value \( f_{\text{PBH}} \leq 10^{-6} \) is two orders of magnitude higher than the upper PBH abundance \( \leq 10^{-8} \) in Ricotti et al. (2008) for \( M_{\text{PBH}} > 10^3 \odot \), but much lower than the lower-mass PBH abundance constraint in Ricotti et al. (2008). However, the massive PBH abundance constraint in Ricotti et al. (2008) is made by the Compton y-parameter estimate at \( z_{\text{rec}} < z \ll z_{\text{ref}} \) without local UCMH trapping, while we consider two-step accretion, first by host UCMHs and then by centre PBHs, as mentioned in Section 2.2.3. X-rays can locally heat the accreted gas inside UCMHs but not the entire cosmic gas at high redshift.

The abundance of IGM molecular hydrogen \( f_{\text{H}_2} \) in various UCMH radiation models is shown in Fig. 6 in this section. We see that when UCMH energy injection can be neglected, this fraction returns to \( f_{\text{H}_2} \sim 10^{-6} \), which is consistent with the standard result (e.g. Galli & Palla 1998 and references therein). The upper bound of enhanced \( f_{\text{H}_2} \) is \( f_{\text{H}_2} \sim 10^{-3} \), due either to the allowed UCMH dark matter annihilation constrained by the CMB optical depth or the X-ray emission at \( z_{\text{min}} \leq 100 \).

**Figure 5.** Ratio of Jeans mass to ‘CMB mass’ (left panels) and evolution of Thomson scattering depth (right panel) for the starting injection redshift \( z_{\text{in}} = 100 \) (upper panels) and 20 (low panels). The lines are as the same in Fig. 4.

### 4 FIRST STRUCTURES

In the hierarchical cold dark matter (CDM) scenario, the first cosmological objects are dark matter haloes, which are formed by gravitational instability from the scale-free density fluctuations (Diemand, Moore & Stadel 2005; Green, Hofmann & Schwarz 2005; Yoshida 2009). The formation of the first baryonic objects depends on the detailed gas dynamical processes. The first baryonic objects can successfully collapse and form inside dark haloes when the cooling time-scale for dissipating the kinetic energy is much shorter than the Hubble time. Tegmark et al. (1997) showed that the formation of baryonic structure crucially depends on the abundance of molecular hydrogen \( f_{\text{H}_2} \). Biermann & Kusenko (2006) considered that photons emitted by dark matter annihilation or decay inside a halo can boost the production of \( \text{H}_2 \), and may favour the formation of the first structure. On the other hand a different conclusion, that dark matter annihilation or decay can slightly delay the first baryonic structure formation, was given by Ripamonti et al. (2007b). As they discussed, the higher central density in the baryonic cloud without dark matter energy injection could compensate for the lower abundance of \( \text{H}_2 \) and still lead to the fastest cooling. Nevertheless, regardless of the promotion or suppression caused by halo extended dark matter annihilation or decay, such effects are very small.

UCMHs can also be captured by the first dark matter objects; in this case, the UCMH radiation can be much brighter than that from the extended large dark halo, even if the fraction of UCMHs in the dark halo is tiny. In this section we focus on dark halo structures with mass \( \sim 10^6 \odot \) (or \( 10^7 \odot \)) in the PBH heating case, as they favour later first star formation (Broom et al. 2009; Yoshida 2009). For a typical \( 10^5 \odot \) halo in the early Universe, \( \sim 10^{20} (10^{18}) \) or \( \sim 10^8 (10^6) \) UCMHs within this halo can be expected for an initial UCMH fraction \( f_{\text{UCMH}} \sim 10^{-4} (10^{-6}) \), if the seeds of UCMHs are generated in electroweak or quantum chromodynamic (QCD) phase transitions respectively (Scott & Sivertsson 2009). In these cases, UCMH emission can also provide a new type of radiation background in the dark halo. However, the uniform UCMH distribution treatment will break down if the number of massive
UCMHs in a halo is less than \( \sim 10 \). This happens for the massive UCMH case \( f_{\text{UCMH}} M_{\text{DM}} / m_\chi \lesssim 10 \). In this section we study the effects of UCMH radiation on the first structure formation and evolution with \( f_{\text{UCMH}} M_{\text{DM}} \gg m_\chi \); the case of ‘only several luminous UCMHs’ will be discussed separately in Section 5.4 as a supplement. As the virial temperature of \( \sim 10^6 M_\odot \) haloes is less than the threshold for atomic hydrogen line cooling, these haloes are often referred as ‘minihaloes’ in the literature. However, for clarity, in this paper we call the first dark matter structures (cosmological) dark matter haloes or dark haloes, which should not be confused with UCMHs.

### 4.1 Dark matter annihilation

The profiles of cosmological dark matter haloes are chosen before our calculation. The equations for the halo profile are listed in the Appendix. Before the formation of the first stars, the energy injection inside a large dark matter halo is mainly contributed by local emission from UCMHs in the halo, local annihilation or decay of the extended dark matter within the halo, and the outside radiation background that is injected into the halo. We check the total energy produced by dark matter annihilation with a dark halo, i.e., \( L_{\text{halo}} = L_{\text{UCMH}} + L_{\text{ext}} \), with \( L_{\text{UCMH}} \) and \( L_{\text{ext}} \) being the annihilation luminosity from UCMHs and the extended dark matter in this halo. The typical ratio \( L_{\text{UCMH}} / L_{\text{ext}} \) is demonstrated in Fig. 7, where we adopt an isothermal dark halo model and \( f_{\text{UCMH}} = 10^{-6} \). We find that both \( L_{\text{UCMH}} \) and \( L_{\text{ext}} \) are proportional to the total halo mass \( M_{\text{DM}} \), so \( L_{\text{UCMH}} / L_{\text{ext}} \) is independent of \( M_{\text{DM}} \). In Fig. 7, \( L_{\text{UCMH}} \) is mostly dominant in the halo, except for large \( z_{\text{vir}} \) with light dark matter particles (e.g., \( < 10 \text{ GeV}/f_{\text{UCMH}} = 10^{-6} \) and \( z_{\text{vir}} = 100 \); much lighter dark matter particles are required for more abundant UCMHs or smaller \( z_{\text{vir}} \)). Therefore, similarly to the IGM environment, the energy injection mechanism inside a dark halo that contains UCMHs can be very different from the no-UCMH case. We focus on the ionization and heating inside the dark matter halo. Also, we mention that the results for \( L_{\text{UCMH}} / L_{\text{ext}} \) in Navarro, Frenk & White (1996) haloes are very similar to the isothermal haloes in Fig. 7.

The simple criterion of baryonic matter filling in a dark halo is approximately taken as the IGM gas temperature being cooler than the virial temperature of the halo, \( T_{\text{m}} < T_{\text{vir}} \), otherwise the gas pressure prevents the gas from collapsing. We do not include the temperature cooling-slope criterion as in Tegmark et al. (1997) to study further baryon cooling and collapse in the dark halo. Fig. 8 gives the dark matter halo mass for the critical case \( T_{\text{m}} = T_{\text{vir}} \) with UCMH annihilation heating the ambient IGM gas. Generally, more massive dark halo mass is needed for gas filling into dark haloes and forming baryonic structure in the haloes. If \( T_{\text{m}} = T_{\text{CMB}} \), the minimum halo mass for gas infilling is \( 6.1 \times 10^3 M_\odot \), while the minimum halo mass increases significantly for \( T_{\text{m}} \gg T_{\text{CMB}} \). In this sense, the formation of the first baryonic objects will obviously be suppressed in small dark haloes located in host IGM gas. However, in the pure UCMH annihilation case without PBH radiation, we expect that \( \sim 10^6 M_\odot \) dark haloes attract baryons at \( z > 10 \) in most cases, except for \( f_{\text{UCMH}} m_\chi^{-1} \gtrsim 10^{-2} \).

### Figure 7

Ratio between \( L_{\text{UCMH}} \) and \( L_{\text{ext}} \) for \( m_\chi c^2 =1 \text{ GeV} \) (solid lines), 10 GeV (dashed lines), 100 GeV (dotted lines) and an isothermal extended dark matter halo. Thick and thin lines correspond to virial redshifts \( z_{\text{vir}} = 20 \) and 100 respectively. A UCMH fraction of \( f_{\text{UCMH}} = 10^{-6} \) is adopted. The case of \( L_{\text{UCMH}} = L_{\text{ext}} \) is marked as the horizontal thin-dashed line.

### Figure 8

The minimal dark halo mass for \( T_{\text{m}} = T_{\text{vir}} \) for an IGM heated by UCMHs. The shaded area is for the case in which the minimal halo mass for gas collapse becomes \( M_{\text{DM}} < 6.1 \times 10^3 M_\odot \), due to a lower ambient gas temperature compared with the CMB temperature. The lines from \( m_\chi c^2 = 100 \text{ MeV–}500 \text{ GeV} \) are as in Fig. 3.
The annihilation energy deposited in a dark halo can be linearly divided into two parts: the energy from the background where cosmic UCMH annihilation occurs, $\epsilon_{\text{bgd}}(z)$, and that within the local dark halo, $\epsilon_{\text{loc}}(z)$. The term $\epsilon_{\text{bgd}}(z)$ is obtained by equation (39). The local energy $\epsilon_{\text{loc}}(z)$ is a function of position inside the halo. We focus on energy deposition at the centre of the halo. The contribution by the extended dark matter in an isothermal halo is

$$
\epsilon_{\text{loc}, \text{iso}}(z) = \frac{\langle \sigma v \rangle c^2 (1 - f_z)}{2 \mu_p m_x f_x} \int_0^{R_{\text{vir}}} \rho^2 \, dr,
$$

where $f_x = \Omega_{\text{DM}}/\Omega_{\text{M}} \approx 0.833$. On the other hand, the local energy deposited by UCMH annihilation depends on the UCMH distribution inside the halo. If we assume that the UCMH number density is uniformly distributed depending on the halo mass density, i.e. $d\sigma_{\text{UCMH}}(r)/dM_{\text{DM}}(r) \propto \text{constant}$ (for a relevant distribution simulation see Sandick et al. 2011), we have

$$
\epsilon_{\text{UCMH}}(z) \simeq \frac{L_{\text{UCMH}}\sigma_{\text{tot}}(E_x)(1 - f_z)}{\mu_p m_x M_{\text{DM}}} \int_0^{R_{\text{vir}}} \rho^2 \, dr
$$

where $f_x = \omega_{\text{DM}}/\omega_{\text{M}} \approx 0.83$. However, after a dramatic temperature increase during the virializing process, $z \sim z_{\text{vir}}$, $H_2$ cooling becomes the main process to cool the denser gas at $z < z_{\text{vir}}$. The peak temperature during virializing is around $\sim 1000$–$2000$ K. We find that higher $H_2$ abundance, which is caused by higher UCMH annihilation, gives a lower gas temperature after virialization for small $z_{\text{vir}}$ ($\sim 20$) but a higher temperature for large $z_{\text{vir}}$ ($\sim 100$). This result is between those in Biermann & Kusenko (2006), who considered that the effects of sterile neutrino decay can favour structure formation, and Ripamonti et al. (2007b), who showed that dark matter annihilation will slightly delay structure formation. The main reason for our difference from Ripamonti et al. (2007b) for $z_{\text{vir}} \ll 100$ is that we take the baryon gas density to be proportional to the halo density $n_b \propto \rho$ as in Tegmark et al. (1997), therefore more molecular gas due to stronger heating just means more efficient cooling. A more elaborate result can be made by adding more detailed gas dynamics and energy transfer including the UCMH radiation within the halo (Tegmark et al. 1997; Ripamonti et al. 2007b). However, such a new calculation should not change the fact that UCMH radiation, as well as extended dark halo annihilation, cannot change the gas temperature in the halo in an obvious way after virialization.

Fig. 9 shows the gas evolution at the centre of a $10^9 M_\odot$ isothermal halo virializing at $z_{\text{vir}} = 20$ or $z_{\text{vir}} = 100$. We also show the protohalo stage at $z > z_{\text{vir}}$. More energetic annihilation due to larger $f_{\text{UCMH}}$ or lower $m_x$ gives higher $x_{\text{H}_2}$ and $f_{\text{L}_2}$. The UCMH annihilation gives a significant impact on $T_m$ before virialization in the protohalo stage; this is because the cooling and heating mechanisms are different before and after virialization. The change of $T_m$ before virialization is mainly due to heating by background UCMH annihilation, which is completely dominant over extended dark matter annihilation; thus, brighter UCMH annihilation luminosity gives a higher gas temperature at $z > z_{\text{vir}}$. However, after a dramatic temperature increase during the virializing process, $z \sim z_{\text{vir}}$, $H_2$ cooling becomes the main process to cool the denser gas at $z < z_{\text{vir}}$. The peak temperature during virializing is around $\sim 1000$–$2000$ K. We find that higher $H_2$ abundance, which is caused by UCMH annihilation, gives a lower gas temperature after virialization for small $z_{\text{vir}}$ ($\sim 20$) but a higher temperature for large $z_{\text{vir}}$ ($\sim 100$). This result is between those in Biermann & Kusenko (2006), who considered that the effects of sterile neutrino decay can favour structure formation, and Ripamonti et al. (2007b), who showed that dark matter annihilation will slightly delay structure formation. The main reason for our difference from Ripamonti et al. (2007b) for $z_{\text{vir}} \ll 100$ is that we take the baryon gas density to be proportional to the halo density $n_b \propto \rho$ as in Tegmark et al. (1997), therefore more molecular gas due to stronger heating just means more efficient cooling. A more elaborate result can be made by adding more detailed gas dynamics and energy transfer including the UCMH radiation within the halo (Tegmark et al. 1997; Ripamonti et al. 2007b). However, such a new calculation should not change the fact that UCMH radiation, as well as extended dark halo annihilation, cannot change the gas temperature in the halo in an obvious way after virialization.
The minimal dark halo mass at are produced < \( \z_{\text{in}} = 100 \) (dark lines), 50 (red lines) and 20 (blue lines). The characteristic PBH radiation is \( E_X = 10^3 \text{keV} \). The shaded area is for the case in which the minimal halo mass for gas collapse becomes \( M_{\text{DM}} < 6.1 \times 10^7 M_\odot \) due to a lower ambient gas temperature compared with the CMB temperature.

Note that the temperature \( T_m \) in the halo only changes by a factor of \( \sim 3 \) for a several orders of magnitude change to the UCMH annihilation luminosity inside the halo. Therefore, we cannot expect that the first baryonic structure formation can be obviously promoted or suppressed. After virialization, the effects of dark matter annihilation are always secondary compared to the H\(_2\) cooling mechanism.

4.2 Gas accretion on to PBHs and X-ray emission

If we consider gas accretion on to PBHs, and take the PBH fraction as \( \eta_{\text{PBH}} > 10^{-11} \) (dash–dotted lines), \( 10^{-8} \) (solid lines), \( 10^{-9} \) (dashed lines), \( 10^{-10} \) (dotted lines) and \( \z_{\text{in}} = 100 \) (dark lines), 50 (red lines) and 20 (blue lines). The characteristic PBH radiation is \( E_X = 10^3 \text{keV} \). The shaded area is for the case in which the minimal halo mass for gas collapse becomes \( M_{\text{DM}} < 6.1 \times 10^7 M_\odot \) due to a lower ambient gas temperature compared with the CMB temperature.

The evolution of the baryonic structure inside a \( 10^7 M_\odot \) isothermal dark halo, as an example, is shown in Fig. 11. We consider two models: the number density of PBH-host UCMHs being uniformly distributed per halo mass, as mentioned in Section 4.1, and also uniformly distributed UCMHs per halo volume inside the halo, i.e. \( dN_{\text{UCMH}} / dV_{\text{halo}} \propto \) constant. Different UCMH distributions inside the halo will give different baryonic evolution. The local energy deposition for the UCMHs with uniform distribution per halo volume is written as

\[
\epsilon_{\text{loc, acc}}(z) \approx \frac{L_{\text{acc}} \sigma_{\text{tot}} E_X \rho_{\text{core}} (1 - f_s)}{\mu m_p f_s V_{\text{halo}}} \left[ 2R_{\text{core}} - \frac{R^2_{\text{core}}}{R_v} \right],
\]

where \( L_{\text{acc}} \) is the total X-ray luminosity due to gas accretion on to PBHs inside the first baryonic objects. The main results in Fig. 11 are very similar to those of dark matter annihilation in Fig. 9. For low virialization redshift, \( z_{\text{vir}} = 20 \), the gas temperature is cooler for higher X-ray luminosity but \( T_m \) inside the halo only changes by a factor of 2–4 for a four orders of magnitude change in X-ray luminosity. A more obvious change of chemical quantities (\( x_{\text{ion}}, f_{\text{H}_2} \)) than the temperature change occurs for different X-ray luminosity within the halo. This means that the effects of X-ray emission from PBHs on the gas evolution inside a halo are very small. PBHs that are uniformly distributed per halo volume give a slightly cooler gas within the halo, as well as lower \( x_{\text{ion}} \) and \( f_{\text{H}_2} \) than uniformly distributed UCMHs per halo mass. That is because the latter model makes the average distance of UCMHs closer to the halo centre and this yields stronger effects to change the gas properties at the centre.

In summary, UCMH radiation including both annihilation and PBH gas accretion enhances the baryon chemical quantities such as \( x_{\text{ion}} \) and \( f_{\text{H}_2} \) inside dark matter haloes above the minimal halo mass for \( T_m = T_{\text{vir}} \), but the impact of UCMH radiation on the temperature of the first baryonic objects is small (by a factor of several), which shows that the change in first baryonic structure formation due to UCMH radiation is less important than the \( \text{H}_2 \) cooling and dark halo virialization time. However, the new chemical conditions provided by UCMH radiation can be more important in affecting later gas collapse and the first star formation after first baryonic object formation, because more abundant \( \text{H}_2 \) and electrons acting as the cooling agents can cool the gas more efficiently during the gas collapse process and provide a lower fragmentation mass scale and first star mass (e.g. Stacy & Bromm 2007). As we mainly focus on the first baryonic structure formation and evolution, the calculation of the first star formation due to the changed gas chemical components should be investigated in more detail in future.

5 DISCUSSION

5.1 Status of UCMH radiation in reionization, other sources

A variety of cosmological sources can reionize and heat the IGM at different redshifts before \( z \simeq 6 \). So far we have shown that UCMHs, even if they merely occupy a tiny fraction of the total dark matter mass, provide a new gamma-ray background for gas heating and ionization. Also, the X-ray emission from accreting PBHs could change the IGM gas evolution history dramatically after \( z_{\text{reion}} \ll 1000 \), where the value of \( z_{\text{reion}} \) depends on the masses of PBH, host UCMHs and dark matter particles. Furthermore, we investigate how both dark matter annihilation and X-ray emission from UCMHs can dominate over the annihilation of extended dark matter haloes. Therefore UCMHs are also an important energy source in dark matter haloes before the first star formation. In this section we briefly review all candidate energy sources during the Universe’s reionization era, \( 10 \leq z < z_{\text{eq}} \). In particular, we emphasize the
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Figure 11. Effects of gas accretion on to PBH inside a $10^7 M_\odot$ isothermal halo with $z_{\text{vir}} = 20$, $z_{\text{bgd}} = 50$, $z_{\text{halo}} = 30$ (left panel) and $z_{\text{UCMH}} = 15$ (right panel). The characteristic $E_X = 10$ keV with $\eta_{f_{\text{PBH}}} = 10^{-7}$ (solid lines), $\eta_{f_{\text{PBH}}} = 10^{-8}$ (dashed lines), $\eta_{f_{\text{PBH}}} = 10^{-9}$ (dotted lines) and $\eta_{f_{\text{PBH}}} = 10^{-10}$ (dash–dotted lines). We choose a more massive halo of $10^7 M_\odot$ because more massive haloes favour structure formation in this case. In the left panel we take $dn_{\text{UCMH}}/dM_{\text{halo}}$ uniformly and in the right panel $dn_{\text{UCMH}}/dV_{\text{halo}}$ uniformly.

importance of UCMH radiation among all these sources at different times.

In this paper we focus on heating and ionization processes after the last scattering epoch $z \sim 1000$, but some interesting effects can be produced by primordial energy sources at earlier times $z > 1000$. For example, cosmic gas heating and CMB spectral distortion at $z_{\text{rec}} < z < z_{\text{eq}}$ produced by PBH gas accretion can be used to constrain the PBH abundance, so that $f_{\text{PBH}} \leq 10^{-8}$ for $M_{\text{PBH}} \geq 10^3 M_\odot$ in the absence of UCMH annihilation (Ricotti et al. 2008). However, as shown in Section 2.2, UCMH annihilation can be more important than PBH gas accretion in the very early Universe even though UCMHs are just beginning to grow at that time. Hotter cosmic gas heated by dark matter annihilation suppresses the PBH gas accretion to become the dominant source to distort the CMB, and even changes the cosmic recombination process, as shown in Fig. 2. The Compton $y$ parameter is likely to be used to constrain the UCMH abundance based on the annihilation scenario in future work.

The influence of dark matter annihilation or decay at $z \leq 1000$ on the IGM during the reionization era has been discussed by many authors (Chen & Kamionkowski 2004; Hansen & Haiman 2004; Pierpaoli 2004; Mapelli & Ferrara 2005; Padmanabhan & Finkbeiner 2005; Zhang et al. 2006; Mapelli et al. 2006; Ripamonti 2007; Chluba 2010; Yuan et al. 2010). Some authors also use observational data to constrain the cross-section of dark matter interaction (Cirelli et al. 2009; Galli et al. 2009; Slatyer, Padmanabhan & Finkbeiner 2009; Kanzaki, Kawasaki & Nakayama 2010). The basic assumption is that the dark matter distribution is smooth and homogeneous at $z \geq 100$. However, the annihilation power can be strongly increased by UCMHs in the early Universe, as discussed in this paper.

The next commonly suggested source of ionization and heating is the first dark objects (dark haloes), which formed approximately at $z \leq 100$. Dark matter haloes enhanced the overall cosmic dark matter annihilation density due to the dark matter concentration in haloes (Iliev, Scannapieco & Shapiro 2005; Oda, Totani & Nagashima 2005; Ciardi et al. 2006; Chuzhoy 2008; Myers & Nusser 2008; Natarajan & Schwarz 2008; Belikov & Hooper 2009; Natarajan & Schwarz 2009; Natarajan & Schwarz 2010). The mass distribution of dark haloes varies from very low mass ($\sim 10^{-6} M_\odot$) to high mass ($\sim 10^{12} M_\odot$) (Green et al. 2005; Diemand et al. 2005; Hooper et al. 2007), depending on the different damping scales due to different dark matter models (Abazajian, Fuller & Patel 2001; Boehm et al. 2005) as well as the mass–halo function (Press & Schechter 1974; Sheth & Tormen 1999). If UCMHs collapse together with homogeneous dark matter, the UCMH annihilation flux could still dominate over the total annihilation flux within dark matter haloes, at least in massive dark haloes with $f_{\text{UCMH}} M_{\text{DM}} \gg m_h$. However, remember that small dark matter haloes contribute a significant part of the total annihilation rate after structure formation. The profiles of earth-mass dark matter haloes and the gamma-ray flux due to annihilation have also been studied recently (Diemand et al. 2005; Ishiyama, Makino & Ebisuzaki 2010). Similarly to UCMH emission, a large enhancement in the annihilation signal is also expected due to emission from dark matter subhaloes as the remnant of structure formation at $z < 60$. The issue of whether UCMHs or small haloes are more important for ionization and heat after $z \sim 60$ should be investigated in future.
The next ionization sources are accreting PBHs, which are located in their host UCMHs, as mentioned in Section 2.2. PBHs with host UCMHs lead to a faster accretion than naked PBHs, but also absorb the X-ray emission due to baryon accumulation within the UCMH. The PBH accretion can only be more important than UCMH annihilation at $z \lesssim z_{\text{in}} < 1000$ with sufficient abundance and radiation efficiency. Keep in mind that the X-ray emission here is from PBHs, or perhaps PBH–UCMH systems, which play an earlier role than the so-called accreting ‘first black holes (BHs)’, which are the remnants of the first stars at $z \approx 15$.

As mentioned in Section 4, dark matter annihilation or X-ray emission affects baryonic structure formation and evolution. Also, it affects the process of the first star formation. The standard first star formation is carried out at $z \approx 20$ (Abel, Bryan & Norman 2002; Broom et al. 2009), but the first star-forming history can be affected by primordial magnetic fields (Tashiro & Sugiyama 2006) or by extended dark matter annihilation in the halo (so-called ‘dark stars’: see Spolyar, Freese & Gondolo 2008; Spolyar, Freese & Gondolo 2009). Previously it was said that the first stars gave the first light to end the cosmic ‘dark age’, which cannot be true if exotic sources such as dark matter annihilation and accreting PBHs are included.

The next ionization sources are more familiar to us. The first stars emitted UV light and produced ‘ionized bubbles’, which could partially ionize the Universe at $z < 20$ directly or could affect the forthcoming formation of next-generation stars and later galaxy formation (e.g. Haiman & Loeb 1997; Wyithe & Loeb 2003; Shull & Venkatesan 2008; Whalen, Huikestaedt & Minckie 2010). The death of the first stars produced ‘first-generation BHs’, which emitted X-rays and ionized the Universe at $z \approx 15$ or even closer (Cen 2003; Ricotti & Ostriker 2004; Madau et al. 2004; Ricotti et al. 2005; Ripamonti 2007; Thomas & Zaroubi 2008). The reionization process was completed after galaxy formation, as galaxies are generally considered to be the main candidates for reionization of the Universe at $z \approx 6$ (Meiksin 2009, and references therein).

Future work that can be carried out includes studying the heating and ionization processes at $z > 1000$ due to annihilation, comparing the total annihilation rate from small dark haloes ($10^{-6} \, M_\odot < M_{\text{DM}} < f_{\text{UCMH}}^{-1} m_{\chi}$, as mentioned above) with that from UCMHs and distinguishing the impacts of different ionization sources using CMB polarization anisotropies and observational constraints from 21-cm spectra. Actually, CMB polarization and the hydrogen 21-cm line are powerful potential probes of the era of reionization with which to constrain early energy sources. The high multipole of polarization anisotropies may be able to distinguish between UCMHs and small dark structures formed at $z < 100$ and to constrain UCMH and PBH abundances further.

### 5.2 Different UCMH profiles

Remember that in Section 2.1, although several UCMH annihilation rates due to various profiles were given in Fig. 1, we chose the UCMH profile as $\rho \propto r^{-3/4}$ with a cut-off at $\rho_{\text{cut}} \propto (t - t_c)^{-1}$. Such a profile is based on the analytical solution of radial infall on to a central overdensity (Bertschinger 1985). A shallower density profile $\rho(r) \propto r^{-1.5}$, which is given if the central accretor is a black hole, or a steeper profile $\rho(r) \propto r^{-3}$, simulated by Mack et al. (2007), will change the total annihilation luminosity of a UCMH significantly. Fig. 12 gives an example of the different annihilation luminosities due to different density profiles in an entire UCMH. A more concentrated dark matter distribution in a steeper profile leads to a much higher total annihilation rate within the UCMH, because the central region of a UCMH contributes a major part to the total annihilation rate. However, we conclude that the overall cosmic annihilation luminosity density equation (10) will not be changed too much, for two reasons. First of all, the $\rho \propto r^{-3}$ profile usually appears in the outer region of a UCMH, but the change in UCMH density profile at the outer region $r \gg r_{\text{cut}}$ will not dramatically change the total annihilation rate, as the density distribution in the central region is crucial to determine the total annihilation rate. Secondly, the contribution of PBH-host UCMHs (with profile $\rho(r) \propto r^{-2.5}$ near the centre) to the overall cosmic annihilation should be much less important than the initial overdensity-seeded UCMHs ($\rho(r) \propto r^{-2.25}$), due to both their shallower inner profile and their much lower abundance $f_{\text{PBH}} \ll f_{\text{UCMH}}$. The annihilation luminosity should still be taken into account if the dark matter particle inner trajectory is highly eccentric with a much closer pericentre than the cut-off radius, as in equation (4), but it is still lower than the luminosity of overdensity-seeded UCMHs as shown in Section 2.1 (Lacki & Beacom 2010).

### 5.3 Extended radiation spectral energy distribution

In Section 3 the products of annihilation of two dark matter particles are assumed to be two gamma-ray photons, both with energy $m_{\chi} c^2$, and X-ray photons emitted from PBH-host UCMHs are characterized by a single energy reflecting the centre PBH mass, $E_X \approx 3 \, \text{keV} (M_{\text{PBH}}/M_\odot)^{-1/4}$. This simplified treatment with a monochromatic (i.e. $\delta$-function) spectrum of local UCMH emissivity is a good approach to demonstrate crucial UCMH effects on the IGM evolution dependent on the most crucial parameters such as UCMH and PBH abundances $f_{\text{UCMH}}$ and $f_{\text{PBH}}$, as well as dark matter particle mass $m_{\chi}$. In this section we will experiment with other extended photon spectral energy distributions (SEDs) for UCMH radiation, rather than a $\delta$-function, and will study the dependence of UCMH...
emissivity on the SED. We will see that more elaborate considerations can quantitatively change the ionization results, but will not qualitatively change the basic conclusions regarding UCMH radiation. We first adopt a power-law spectrum \( F(E, z) \propto E^{-\alpha} \) and then discuss another spectrum \( F(E, z) \propto E^\beta \exp(-dE) \); these are the two most commonly used SEDs for annihilation and BH X-ray emission.

First of all we give an analytic calculation for UCMH radiation with a locally monochromatic spectrum, based on an approximation that the optical depth described in equation (38) can be neglected, \( \tau < 1 \). This condition is applied to the spectrum \( E_i(E_{\chi}) \gg 10 \text{ keV} \). Under the approximation \( \tau \ll 1 \), equation (37) in Section 3 can be simplified as

\[
\frac{dn}{dE} \simeq \frac{1}{E_0 E_{\gamma}} \left( \frac{E_{\gamma}}{E_0} \right)^{(1/2)-z} \frac{l(z)}{(1+z)^{3/2}},
\]

where the emitted photon number \( n \) and luminosity \( l(z) \) can be applied to both annihilation \([n_{ann}, l_{ann}(z)]\) and X-ray luminosity \([n_X, l_{X}(z)]\), \( E_0 \) for the characteristic energy \( E_{\gamma} \) or \( E_X \) respectively, and \( l(z) \propto (1+z)^{\xi} \) with \( \xi \approx 4 \) for annihilation and \( \xi \approx 3 \) for X-ray emission. We take the interaction cross-section between photons and the IGM gas as \( \sigma_{\text{tot}}(E) \propto E^{-\delta} \). This condition is applied to both annihilation \([n_{ann}, l_{ann}(z)]\) and X-ray luminosity \([n_X, l_{X}(z)]\), \( E_0 \) for the characteristic energy \( E_{\gamma} \) or \( E_X \) respectively, and \( l(z) \propto (1+z)^{\xi} \) with \( \xi \approx 4 \) for annihilation and \( \xi \approx 3 \) for X-ray emission. When \( \xi \approx 4 \), we only track \( l(z) \) based on different SEDs.

The first extended SED is \( F(E, z) \propto E^{-1} \) for \( E_1 \leq E \leq E_2 \), which gives a spectral number density for \( \tau \approx 0 \) of

\[
\frac{dn}{dE} \simeq \frac{c}{H_0 \Omega_{\text{m}}} \left( \frac{2}{13 - 2z} \right) \frac{l(z)}{\ln \Lambda} (1+z)^{-3/2} E_{\gamma}^{-2},
\]

where the subscript \( \delta \) is marked for a monochromatic SED as in equation (40). For high-energy gamma-ray photons we can take \( k \approx 2 \) for the Klein–Nishina cross-section, but for X-ray photons we take \( k \approx 0 \) for \( E_X \gg 1 \text{ keV} \). Another thing we mention is that ionization rate \( \dot{\zeta}(z) \propto \epsilon(z) \) for \( x_{\text{ion}} \ll 1 \), so we only track \( \epsilon(z) \) based on different SEDs.

If we do not focus on the linear changed factors such as \((a-1)/a + (1/2) - \zeta \), in most cases the power-law spectrum will significantly increase the energy density \( \epsilon(z) \) as well as the ionization rate \( \dot{\zeta}(z) \), for example in some cases by a factor of \( (E_0/E_{\gamma})^2 \). A blackbody-like or multicomponent spectral distribution \( F(E, z) \propto E^{\beta} \exp(-dE) \) with a peak \( E_0 \) and \( c \gg 0 \) can be approximately written as \( F(E, z) \propto x^\delta \) for \( x < 1 \) and \( F(E, z) \propto \exp(-dE) \) for \( x > 1 \), where \( x = E/E_0 \). In this case we find the number-density spectrum as

\[
\frac{dn}{dE} \simeq \frac{cA}{H_0 \Omega_{\text{m}}} \left( \frac{E_0}{E_{\gamma}} \right)^{b} \left( 1 + \frac{b - \zeta}{E_{\gamma} / E_X} \right),
\]

where \( A \) is an algebraic factor \( A = [(1 + b)^{-1} + d^{-1} - 1] \) and \( 0 < f(d) < 1 \) can be obtained numerically, which is not important for the following discussion. The terms in the number-density spectrum are proportional to \( E_{\gamma}^0 \) or \( E_{\gamma}^{1/2-\zeta} \). According to equation (58), the integrated energy density \( \epsilon(z) \) is expected to be similar to equation (52) for \( b > k - 1 \). Otherwise the enhancement should be \( \epsilon(z) \propto \epsilon_s(z)(E_d/E_{\gamma})^{b_k} \). A steep \( x^\delta \) with \( c > 1 \) for annihilation \( k \approx 2 \) and all the blackbody-like SEDs for X-rays \( k \approx 0 \) are more or less similar to the \( \delta \)-function SED at \( E_0 \).

Now we study the case of \( E_X \ll 10 \text{ keV} \) for X-ray emission from very massive PBHs. Photon absorption in the IGM is important for \( E_X \sim 1 \text{ keV} \). The mean free path of X-ray photons undergoing redshift change \( \Delta z \) can be written as

\[
\tau \sim \frac{\Delta z}{H_0 \Omega_{\text{m}}} \left( E_{\gamma} / E_X \right)^{1/2} \epsilon_s(1 \text{ keV}) \left( E_X / 1 \text{ keV} \right)^{-k},
\]

or we have

\[
\Delta z \sim 0.40(1 + \zeta)^{-1/2} \left( E_X / 1 \text{ keV} \right)^{k},
\]
where \( k \approx 3.3 \) for the photionization cross-section. Similarly to the former calculation, the energy density for a \( \delta \)-function spectrum is

\[
\epsilon_\delta(z) \sim \frac{0.40}{H_0 \sqrt{\Omega_{\text{ME}}} \Lambda} \int_0^z n_A(E_0) \left( \frac{E_0}{1\,\text{keV}} \right)^k \sigma_{\text{ion}}(E_0) \, dz.
\]  

Note that we now have a shallower density evolution \( \epsilon_\delta(z) \propto (1 + z)^{3/2} \) compared with the transparent propagation for the case of equation (54), \( \epsilon_\delta(z) \propto (1 + z)^3 \). For a power-law SED, \( F(E, z) \propto E^{-\alpha} \), we obtain the energy density

\[
\epsilon(z) \sim \frac{0.40 c}{H_0 \sqrt{\Omega_{\text{ME}}} \Lambda} \int_0^z n_A \sigma_{\text{ion}}(E_0) \left( \frac{E_0}{1\,\text{keV}} \right)^{k-1} \alpha^{-1}.
\]  

Therefore the energy density is enhanced for \( E_1 > E_0 \), but the enhancement is less significant compared with that of equation (57) for the same spectral index \( \alpha \).

In brief summary, the strength of UCMH radiation depends on its extended spectral energy distribution, which will increase the properties of IGM ionization \( I(z) \) and heating \( \epsilon(z) \). Compared with the basic results with a monochromatic spectrum \( E_x \) for annihilation or \( E_X \) for X-ray emission, a power-law spectrum \( \propto E^{-\alpha} \) with \( \alpha > 0 \) can increase the energy density effectively, but a blackbody-like spectrum is more like the monochromatic SED case. For local heating of \( E_X \approx 1\,\text{keV} \), the heating increases are less significant than for the transparent propagation case \( \tau \ll 1 \). Also, a power-law spectrum changes \( \epsilon(z) \) and \( I(z) \) more significantly for annihilation than for X-ray radiation.

### 5.4 More massive UCMHs inside the first dark haloes

In Section 4 we assume that the number of UCMHs inside a dark halo is so huge that UCMHs in the halo are uniformly distributed per halo mass \( \Delta n_{\text{UCMH}}(dM_{\text{DM}} = \text{const}) \) or per volume \( \Delta n_{\text{UCMH}}(dV_{\text{DM}} = \text{const}) \). This assumption can be invalid if the mass seed of a single UCMH \( m_0 \) is comparable to \( f_{\text{UCMH}}M_{\text{DM}} \). In this case, the position of each UCMH is important to determine the energy deposition within the halo. Usually X-ray emission can be neglected in a massive UCMH due to photon trapping (Section 2.2.3), so we only focus on UCMH annihilation. Equation (49) in Section 4.2 is invalid for massive UCMHs, \( m_0 \sim f_{\text{UCMH}}M_{\text{DM}} \). In this case \( \epsilon_{\text{loc},\text{ucmh}} \) in equation (49) should be calculated as the summation of individual UCMHs:

\[
\epsilon_{\text{loc},\text{ucmh}}(z) \approx \frac{\sigma_{\text{us}}(1 - f_{\text{is}})}{m_p f_X} \sum_i L_{\text{UCMH}}(r_i) \frac{d(r_i)}{4\pi r_i^2}.
\]  

where \( L_{\text{UCMH}} \) and \( r_i \) are the annihilation luminosity and position of the \( i \)th UCMH within the halo.

An extreme case is the one in which only one massive UCMH is inside a dark halo; the energy deposited in the gas at the centre of the halo can then be obtained by equation (63) in Section 4.1. Remember that for uniformly distributed UCMHs with the same total mass as the single UCMH inside a halo we use equation (49). The ratio factor between (63) and (49) is

\[
\text{factor} \approx \left( \frac{R_{\text{core}}}{r_0} \right)^4 \left( \frac{3f_X}{R_{\text{core}}} - 2 \right) \left( 4 - \frac{R_{\text{core}}}{R_{\text{tr}}} \right)^3 -1
\]

\[
\approx \frac{3}{4} \left( \frac{R_{\text{core}}}{r_0} \right)^4,
\]  

where \( r_0 \) is the location of the single massive UCMH from the halo centre. If we take \( R_{\text{core}} = \xi R_{\text{tr}} \) with \( \xi \ll 1 \), the factor can be approximately written as \((3/4\xi)(R_{\text{core}}/r_0)^4\). For a single UCMH located close to the isothermal case \( r_0 \approx R_{\text{core}} \), the factor (64) is \((3/4\xi)\). For a single UCMH located at the turnaround radius \( r_0 \approx R_{\text{tr}} \), equation (64) becomes \((3\xi^3)^{1/4}\). Taking a typical value \( \xi \approx 0.1 \), the energy deposition \( \epsilon_{\text{ucmh}} \) contributed by a single UCMH varies from \( \sim 10 \) higher to \( \sim 10^{-1} \) lower compared with that contributed by the same total mass but uniformly distributed small UCMHs. We also check the importance of a single UCMH’s position on the gas temperature inside a halo, but we see that the effect on temperature can be neglected.

On the other hand, the energy deposition by a single UCMH compared with a volume of uniformly distributed small UCMHs with the same total mass can be amplified by a factor

\[
\text{factor} \approx \frac{1}{6} \left( \frac{R_{\text{core}}}{r_0} \right)^4 \left( \frac{R_{\text{tr}}}{R_{\text{core}}} \right)^3 \left( 1 - \frac{R_{\text{core}}}{R_{\text{tr}}} \right)
\]

\[
\approx \frac{1}{6\xi^3} \left( \frac{R_{\text{core}}}{r_0} \right)^4,
\]  

which gives an enhancement factor of about \( \sim 1/(6\xi^3) \sim 1/\xi^2 \) for \( r_0 = R_{\text{core}} \) and a weaker factor \( \xi/6 \) for \( r_0 = R_{\text{tr}} \). These results show that a massive single UCMH is always more important for energy deposition compared with a volume of uniformly distributed UCMHs.

Therefore we conclude that different UCMH distributions with the same total mass but different individual mass can change the energy deposition and structure evolution inside a halo. A more concentrated distribution towards the halo centre or a more closely located single halo near the centre gives a more significant effect on gas ionization and heating at the halo centre.
6 CONCLUSIONS

Ultracompact minihaloes (UCMHs) have been proposed as primordial dark matter structures that formed by dark matter accretion on to an initial overdensity or primordial black holes (PBHs) after matter–radiation equality, $\zeta_\text{eq} \approx 3100$ (Mack et al. 2007; Ricotti & Gould 2009). The key difference between UCMHs and the first dark halo structures is that UCMHs are seeded by primordial density perturbations produced in the very early Universe, such as phase-transition epochs $(10^{-3} \leq \delta \leq 0.3$ for the initial overdensity or $\delta > 0.3$ for PBHs), so they can grow short after $z \sim \zeta_\text{eq}$. The radiation from UCMHs in the early Universe includes dark matter annihilation from all UCMHs and X-ray emission from gas accretion on to PBHs. In this paper we investigate the influence of UCMH radiation on the early ionization and thermal history of the IGM and the subsequent evolution of the first massive baryonic objects. Our conclusions are as follows.

(1) UCMH annihilation can completely dominate over homogeneous dark matter background annihilation and can provide a new gamma-ray background even for a tiny UCMH fraction, $f_{\text{UCMH}} = \Omega(\zeta_\text{eq})^2/\Omega_{\text{DM}} \sim 2.2 \times 10^{-19} m_{\chi_\text{UCMH}}^{-2/3} (1 + z)^2$ with $m_{\chi_{\text{UCMH}}} = m_\chi \sqrt{e/100}\text{GeV}$. We conclude that the influence of dark matter annihilation on IGM evolution can be significantly enhanced when we include UCMHs in addition to the homogeneous dark matter background. In most cases, UCMH annihilation has been the dominant source of ionization and gas heating since the matter–radiation equality epoch, until X-ray emission from PBHs or large-scale structure formation became important at $z \lesssim 100$.

(2) The impact of UCMH annihilation on the IGM can be approximately estimated by the quantity $m_{\chi_{\text{UCMH}}}^{-1} f_{\chi_{\text{UCMH}}}$. The threshold for UCMH abundance $f_{\chi_{\text{UCMH}}} \geq 10^{-6}$, after the matter–radiation equality epoch, the IGM ionization fraction $x_{\text{ion}}$ can be increased from $x_{\text{ion}} \sim 10^{-4}$ in the absence of any energy injection to an upper bound $x_{\text{ion}} \lesssim 0.1$, and the IGM temperature from the adiabatic cooling $T_m \propto (1 + z)^2$ to a maximum value $T_m \sim 5000\text{K}$ for the upper-bound case $m_{\chi_{\text{UCMH}}}^{-1} f_{\chi_{\text{UCMH}}} \sim 10^{-2}$, which is constrained using the CMB optical depth at late times $\tau < 30$. UCMH annihilation is able to increase significantly the Thomson optical depth $\tau \geq 0.1$ in the early Universe $z \gg 30$, which is unrelated to the measured CMB optical depth at $z < 30$. The UCMH annihilation luminosity is based on the UCMH profiles, where we take $\rho \propto r^{-2.25}$ from the literature (Bertschinger 1985). Steeper (shallower) profiles decrease (increase) the allowed upper limit of $f_{\chi_{\text{UCMH}}}$, but the variations in the overall IGM chemical and thermal quantities should not be changed too much, because the fraction of UCMHs with a profile $\rho \propto r^{-1.5}$ as PBH hosts is very small, and $\rho \propto r^{-3}$ occurs only in the halo outskirts $r \gg r_{\text{cut}}$.

(3) Each PBH is located in its host UCMH (Mack et al. 2007). We emphasize that the impact of X-ray emission from PBH-host UCMH systems is limited by the low abundance of PBHs ($f_{\text{PBH}} \ll f_{\chi_{\text{UCMH}}}$, the average inefficient radiation ($\eta \ll 1$), the photon-trapping effect by accreted baryons in host UCMHs and outflows produced by rapid accretion feedback. Sufficiently massive host UCMHs can accrete and thermalize infalling baryons, which are accumulated inside UCMHs with a mass fraction $f_{\text{gas}} > 10^{-3}$, and trap X-rays from accreting PBHs until a critical redshift $z_{\text{cut}} \sim 32(5m_{\text{PBH}}/m_{\chi_{\text{UCMH}}})^{1/2} m_{\chi_{\text{UCMH}}}^{5/12}$, below which X-rays from super-Eddington accretion flows on to PBHs could escape the surrounding baryon environment in host UCMHs. Although the PBH abundance is $f_{\text{PBH}} \ll f_{\chi_{\text{UCMH}}}$ due to a much higher perturbation threshold for PBH formation, X-ray emission could dominate over UCMH annihilation and become a more promising cosmic energy source of IGM ionization and heating at $z \ll z_{\text{cut}}$ if the PBH abundance is above a threshold of $f_{\text{PBH}}/f_{\chi_{\text{UCMH}}} \sim 3.1 \times 10^{-3} m_{\chi_{\text{UCMH}}}^{1/3} (1 + z)$, which is only allowed beyond the standard Gaussian density perturbation scenario.

(4) As UCMHs are expected to exist in our Galaxy, we expect that UCMHs collapse with the homogeneous dark matter background to form the first large-scale dark matter objects (dark haloes). If this is the case, the dark matter annihilation from UCMHs inside the first dark haloes still dominates over the extended dark matter annihilation background inside the haloes even after halo virialization. UCMH radiation, including both dark matter annihilation and accretion emission, can dramatically suppress the formation of low-mass first baryonic structure, since UCMH radiation heats the IGM and provides a hot ambient gas environment up to $T_m \sim 10^4\text{K}$. The UCMH radiation enhances baryon chemical quantities such as $x_{\text{ion}}$ and $f_{\text{H}_2}$ by orders of magnitude from $x_{\text{ion}} \sim 10^{-6}$ and $f_{\text{H}_2} \sim 10^{-4}$ to the upper bound of $x_{\text{ion}} \sim 10^{-2}$ and $f_{\text{H}_2} \sim 5 \times 10^{-3}$. However, the impact of UCMH radiation on the baryon temperature of the first baryonic objects is very small, which shows that the influence of UCMH radiation on the temperature of the first baryonic objects is small compared with the molecular hydrogen cooling and virialization time $\tau_{\text{vir}}$. However, the more abundant $x_{\text{ion}}$ and $f_{\text{H}_2}$ provided by UCMH radiation decrease the gas temperature in the later gas collapse phase and can produce a lower fragmentation mass scale and lower mass first stars.

Also, we point out that different spectral energy distributions of UCMH radiation also affect the processes of ionizing radiation and heating gas. A more concentrated UCMH distribution within a dark matter halo provides a more promising ionization phenomenon for the gas in the first dark haloes. UCMHs should be distinguished from small dark structures that formed during the structure formation epoch after $z \sim 100$. Future work needs to be done to investigate the importance of small dark matter structure down to Earth mass, compared with UCMH radiation in the early Universe at $z \sim 60$. Also, the CMB polarization anisotropies, 21-cm spectrum and Compton $y$ parameter affected by UCMH radiation also need to be studied further for a better constraint on UCMH abundance.

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Ripamonti (2007) and Ripamonti et al. (2007a). A dark matter halo with a mass \( M_{DM} \) is assumed to be distributed inside a truncation radius \( R_{\text{tr}} \), which is given by

\[
R_{\text{tr}}(z, z_{\text{vir}}) = \begin{cases} 
\left[ \frac{3}{4\pi} \frac{M_{DM}}{\rho_0(z)} \right]^{1/3}, & z \geq z_{\text{ta}} \\
\frac{R_{\text{vir}}}{2} \left[ 2 - \frac{t(z) - t(z_{\text{ta}})}{t(z_{\text{vir}}) - t(z_{\text{ta}})} \right], & z_{\text{vir}} \leq z \leq z_{\text{ta}} \\
R_{\text{vir}}, & z \leq z_{\text{vir}}.
\end{cases}
\]  

(A1)

Here \( z_{\text{vir}} \) is the redshift of halo virialization and \( z_{\text{ta}} \approx 1.5(1 + z_{\text{vir}}) \) − 1 is the turnaround redshift. The dark matter mass inside the halo \( M_{DM} \) is given as a parameter here. Dark matter within the halo is assumed to be distributed uniformly at \( z > z_{\text{ta}} \) with a density evolution

\[
\rho_{\text{DM}}(z, z_{\text{vir}}) = \rho_{\text{DM}}(z_{\text{ta}})^{1.94A/(1 - 0.75A)^2},
\]  

(A2)

where \( A(z) = (1 + z_{\text{vir}})/(1 + z) \). The virial radius \( R_{\text{vir}} \) is given by

\[
R_{\text{vir}} = \frac{1}{2} R_{\text{tr}}(z_{\text{ta}}) = \frac{1}{2} \left[ \frac{3}{4\pi} \frac{M_{DM}}{\rho_0(z_{\text{ta}})} \right]^{1/3}.
\]  

(A3)

For an isothermal halo profile with a core radius \( R_{\text{core}} \) inside \( R_{\text{vir}} \), the dark matter density \( \rho(z) \) is a constant \( \rho_{\text{core}} \) for \( r \leq R_{\text{core}} \), \( \rho(z) \propto r^{-2} \) for \( R_{\text{core}} \leq r \leq R_{\text{vir}} \). The density tends to the background dark matter density for \( r > R_{\text{vir}} \). Here the core radius \( R_{\text{core}} \) is obtained as

\[
R_{\text{core}}(z, z_{\text{vir}}) = \begin{cases} 
R_{\text{tr}}(z), & z \geq z_{\text{ta}} \\
\frac{R_{\text{vir}}}{2} \left[ 2 - (2 - \xi) \frac{t(z) - t(z_{\text{ta}})}{t(z_{\text{vir}}) - t(z_{\text{ta}})} \right], & z_{\text{vir}} \leq z \leq z_{\text{ta}} \\
\xi R_{\text{vir}}, & z \leq z_{\text{vir}}.
\end{cases}
\]  

(A4)

where the coefficient \( \xi \) is introduced as a parameter and \( \rho_{\text{core}} \) can be obtained by integrating the halo mass within \( R_{\text{core}} \) as \( M_{DM} \). The total luminosity produced by the annihilation of dark matter particles distributed within the isothermal halo is

\[
L_{\text{ext,iso}} = \frac{2\pi}{m_f} \left( \frac{\sigma v}{\pi} \right) c^2 R_{\text{core}}^2 \left( R_{\text{tr}} - \frac{2}{3} R_{\text{core}} \right).
\]  

(A5)

On the other hand, the widely used NFW dark matter profile is \( \rho(r) \propto r^{-1}(1 + r/R_{\text{core}})^{-1} \) for \( r \leq R_{\text{vir}} \). Similarly, we can write the complete formula for the density distribution as a function of redshift and the luminosity produced by dark matter annihilation.

The virial temperature of a halo \( T_{\text{vir}} \) is calculated as

\[
T_{\text{vir}} = \frac{\mu m_p G M_{DM}}{2k_B R_{\text{vir}}},
\]  

(A6)

where sometimes the total mass \( M_{\text{halo}} \) is used instead of \( M_{DM} \). However, as we consider whether gas from the background can collapse into the halo, we adopt the pure dark halo mass in order to calculate the initial virial temperature. Therefore we obtain a virial temperature of \( T_{\text{vir}} = 8.2 \times 10^3 \text{K} (1 + z_{\text{vir}}) (M_{DM}/M_\odot)^{2/3} \), or \( T_{\text{vir}} = 380 \text{K} (M_{DM}/10^6 M_\odot)^{2/3} (1 + z_{\text{vir}})/100 \). Baryon-falling can occur when the IGM temperature \( T_m \) around the halo is \( T_m < T_{\text{vir}} \), otherwise the gas pressure will impede gas from falling into the halo and collapsing to form smaller structures. Moreover, the condition \( T_{\text{CMB}} = T_{\text{vir}} \) with \( T_{\text{CMB}} = 2.73 \text{K} (1 + z_{\text{vir}}) \) gives a critical halo mass \( M_{DM} = 6.1 \times 10^3 M_\odot \).

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