Light Dark Matter and Dark Radiation

Jae Ho Heo and C.S. Kim

Department of Physics and IPAP, Yonsei University, Seoul 120-479, Korea

Abstract

The light ($M \leq 20$ MeV) dark matter particles freeze out after neutrino decoupling. If the dark matter particle couples to neutrino or electromagnetic plasma, the late time entropy production by dark matter annihilations can change the neutrino-to-photon temperature ratio, and equally effective number of neutrinos $N_{\text{eff}}$. We study the effect of dark matter annihilations in the thermal equilibrium approximation and non-equilibrium method (freeze-out mechanism), and constrain both results with Planck observations. We demonstrate that the bound of dark matter mass and the possibility of the existence of extra radiation particles are more tightly constrained in the non-equilibrium method.

PACS numbers: 95.35.+d, 98.80.Cq
I. INTRODUCTION

Photons and neutrinos are lightest particles in the Standard Model (SM), and give the radiation energy density at late times of early universe. The SM neutrino species contributes three degrees of freedom, since there are exactly three neutrino mass eigenstates ($\nu_1, \nu_2, \nu_3$), corresponding to the three flavor eigenstates ($\nu_e, \nu_\mu, \nu_\tau$) of the weak interaction. The weak interactions that keep neutrinos in thermal contact with the electromagnetic plasma become ineffective around a second after the Big Bang. Neutrinos decouple at a temperature of order $2 - 3$ MeV before $e^\pm$ pairs annihilate, and thus do not share in the entropy transfer from $e^\pm$ pairs. It causes neutrino temperature to be less than photon temperature later. However, neutrino decoupling is not quite complete when $e^+e^-$ annihilation began, so some of the energy and entropy of photons can transfer to neutrinos. If the dark radiation density is parameterised in terms of the number of effective neutrino species $N_{\text{eff}}$ with the canonical neutrino-to-photon temperature ratio, $N_{\text{eff}}$ increases to slightly more than three neutrino species, leading to $N_{\text{eff}}^{\text{SM}} = 3.046$ [1, 2]. Because the number of effective neutrino species $N_{\text{eff}}$ is precisely predicted in the SM, this can give a robust constraint to any nonstandard physics. For example, new relativistic particles, such as light sterile neutrino [3] or Goldstone boson [4] which has its decoupling temperature less than $100$ MeV, arise in many extensions of the SM, and the existence of them will contribute to the dark radiation energy density. This scenario is, however, strongly excluded at over the $3\sigma$ level in the latest Planck analysis [5], unless photons or electrons (positrons) are heated at later time [6–9].

According to the recent analysis of the cosmic microwave background (CMB) temperature anisotropy by the Planck satellite [5], they found $N_{\text{eff}} = 3.15 \pm 0.23$ (1$\sigma$), consistent with the SM prediction. We should recognize that the Hubble constant ($H_0 = 67.8 \pm 0.9$ kms$^{-1}$Mpc$^{-1}$) inferred by the Planck is in tension at about 2.4$\sigma$ with the direct measurement of $H_0 = 73.8 \pm 2.4$ kms$^{-1}$Mpc$^{-1}$ by HST [10]: larger values of the Hubble constant prefer larger values of $N_{\text{eff}}$. The Planck result must have a difficulty to constrain the dark radiation because of accuracy of astrophysical measurements which have to be combined with. $N_{\text{eff}}$ was not strongly excluded to about the 2$\sigma$ upper limit in their analysis. Additionally, $N_{\text{eff}}$ can be inferred from the big bang nucleosynthesis (BBN) considerations [11] at the time earlier than recombination. The primitive abundance of light elements deduced from astronomical observations can set up the effective number of neu-
trino species if we constrain $N_{\text{eff}}$ as a free parameter. Recently, two groups showed results with a difference from an analysis of $^4\text{He}$ abundance measurements. It was determined in the combination with the D abundance $^{12, 13}$. One group in Ref. $^{14}$ obtained a result which is consistent with the Planck observation, but the other $^{15}$ found a larger value of $N_{\text{eff}} \sim 3.58$. This probe does not have the same resolving power as the Planck satellite. Therefore, we will use Planck results to constrain the dark radiation in this work, and the BBN results are considered in case.

Recently, the light weakly interacting massive particle (WIMP) $^{16-21}$, which can be understood as a dark matter (DM), got some attention, because the existence of a light WIMP can modify the early universe energy and entropy densities, so that it might be able to explain possibly a small difference of $N_{\text{eff}}$ from the prediction or to avoid the tight constraint of experimental measurements for the existence of any other radiation particle. If DM particles couple to the SM particles (neutrinos, or photons and $e^\pm$ pairs) and are sufficiently light ($M \leq 20$ MeV) to annihilate after neutrino decoupling, the late time annihilation of the DM will heat either neutrinos or photons. This will affect neutrino-to-photon temperature ratio and give a contribution to the number of effective neutrino species $N_{\text{eff}}$. This scenario was studied in the equilibrium version for neutrino heating $^7, 8, 22$ and photon heating $^6, 9$. The equilibrium version is, however, a rough approximation, because DM particles are nonrelativistic at freeze-out$^1$. If the DM particles are relativistic at the DM decoupling, they will decouple at the equilibrium concentrations such as the behavior of light or massless neutrinos. Mostly, WIMPs are nonrelativistic at freeze-out due to their weak interaction rate with the thermal background, and the evolution of the number of WIMPs depends upon the annihilation cross section. Here, the Boltzmann equation is applied to the time-evolution of the DM number in spatially homogeneous and isotropic universe. We treat an adiabatic expansion of the universe, so that the total entropy stays constant and the second law of thermodynamics can be applied to the entropy (temperature) evolution of the produced relativistic particles. The parameter that determines the dark radiation energy is the DM mass $M$, so it will be the parameter ($M$) that we will constrain. We start to review the DM number evolution on the expanding universe and the dark radiation

$^1$ We distinguish terminologies, “decoupling” and “freeze-out”, in this paper. The “decoupling” will be used in case that the particle becomes free as maintaining the equilibrium, and the “freeze-out” is for its evolution beyond equilibrium.
\((N_{\text{eff}})\) in equilibrium approximation. Then, we study the out-of-equilibrium light particle production, its entropy (temperature) evolution and its effect on the dark radiation. The possibility of the existence of new light species (equivalent neutrinos) is also investigated.

II. THEORETICAL DETAILS AND NUMERICAL RESULTS

The dark energy density of the universe \(\rho_{\text{DR}}\) is parameterised in terms of the energy density of photons \(\rho_{\gamma}\) and the effective number of neutrinos \(N_{\text{eff}}\) with the neutrino-to-photon temperature ratio of the SM,

\[
\frac{\rho_{\text{DR}}}{\rho_{\gamma}} = \frac{7}{8} N_{\text{eff}} \left(\frac{T_\nu}{T_\gamma}\right)^4_{\text{SM}}.
\]

The factor \(7/8\) is due to the effect of Fermi-Dirac statistics on energy density. The exact formula of \(N_{\text{eff}}\) depends on scenarios (models). Since the temperature will be changed in our scenario, \(N_{\text{eff}}\) can be expressed by

\[
N_{\text{eff}} = N_{\text{eff}}^{\text{SM}} \left(\frac{T_\nu}{T_\gamma}\right)^4_{\text{SM}} \left(\frac{T_\nu}{T_\gamma}\right)^{-4}_{\text{SM}}.
\]

We have assumed that DM particles are nonrelativistic at the time we consider, for example, the time of DM freeze-out or recombination. The BBN imposes limits on \(N_{\text{eff}}\) at the photon temperature around \(1 - 0.1\) MeV, and any additional dark radiation particle is unfavorable from BBN considerations though there is still a small possibility. If a DM particle is relativistic in the BBN era, it becomes a dark radiation particle. We investigate the effect of DM annihilations for the DM mass range of \(0.1 - 20\) MeV. The ratio of neutrino to photon temperatures can be determined by considering entropy conservation, because the total entropy stays constant in adiabatic expansion of the universe. After neutrino decoupling, the primeval plasma would consist of two decoupled components, the electromagnetic and three neutrino ones. The entropy of the neutrino and electromagnetic plasmas are thus separately conserved, and this must serve as an efficient tool for the study of dark radiation here.

There are two independent thermal baths after neutrino decoupling, and we will consider the DM interaction (annihilation)\(^2\) in each thermal bath. Here, the DM particle always interacts with the plasma in thermal bath \(a\), and the thermal bath \(b\) is not relevant to the DM interaction, \(i.e., a, b = \nu\) or \(\gamma\), but \(a \neq b\).

\(^2\) If DM particles are in thermal contact with photons, electrons and neutrinos after neutrino decoupling, it will heat up them equally and the \(N_{\text{eff}}\) will not be changed.
A. Dark Matter Number Evolution

DM particles were in the thermal contact with the rest of the cosmic plasma at high temperatures, but they would experience the freeze-out at a critical temperature. In this case, we should consider the Boltzmann equation for DM number evolution. Since the DM interacts with one of the plasmas, we express the comoving number density about the temperature of the plasma which is not relevant with the DM. This notation is very useful later because one plasma can be always in the thermal equilibrium. If the DM interacts with the plasma in a thermal bath $a$, the evolution equation for the comoving number density $Y$ ($\equiv n_{DM}/s_b$) with respect to the inverse temperature $x_b$ ($\equiv M/T_b$) of the other thermal bath $b$ reads

$$\frac{dY}{dx_b} = -\frac{\langle \sigma v \rangle s_b}{x_b H} (Y^2 - Y_{eq}^2), \quad (3)$$

where $H$ is the Hubble parameter, $s_b$ is the entropy density in the thermal bath $b$ and $Y_{eq}$ ($= n_{eq}/s_b$) is the equilibrium number density. This equation is not meaningful in case that $Y = Y_{eq}$. This becomes a usual fluid equation in the thermal equilibrium. We parameterize the annihilation cross section as $\langle \sigma v \rangle = \sigma_0 x_b^{-n}$ in which $n = 0$ for $s$-wave annihilation and $n = 1$ for $p$-wave annihilation. The above equation can be reduced to

$$\frac{dY}{dx_b} = -\sqrt{\frac{\pi}{45}} m_{pl} M \sigma_0 \left( \frac{g_{ss}^b}{g_s} \right) x_b^{-n-2} (Y^2 - Y_{eq}^2), \quad (4)$$

where $g_s$ and $g_{ss}^b$ are the effective relativistic degrees of freedom for energy density and entropy, and $m_{pl}$ is the Planck mass.

Unfortunately, there is no analytic solution of Eq. (4). Fig. 1 shows the result of numerical solutions for the $s$-wave annihilation into neutrinos (left panel) and $p$-wave annihilation into photons (right panel). The residual annihilations of the DM into photons can distort the CMB spectrum [23, 24]. This effect excludes the DM with masses less than 10 GeV by the $s$-wave annihilation into photons. The effect is negligible for the $p$-wave annihilation which is velocity dependent, and so this bound can be evaded. For the annihilation into neutrinos, we assumed that the same number of neutrinos and antineutrinos of each type was produced in the DM annihilation. The number of the effective relativistic degrees of freedom $g_{ss}^b$ is not relevant with the DM or its production. We could take the value on the SM base. The $g_s$ is taken as a constant $\bar{g}_s$ in an average. The curves were made with the proper values of $\sigma_0 \bar{g}_s^{-1/2}$ in the agreement of the current DM relic density for several
FIG. 1: The comoving number density $Y$ as a function of inverse temperature $x_\gamma (= M/T_\gamma)$ and $x_\nu (= M/T_\nu)$ for a $g = 1$ real scalar (short dash), $g = 2$ Majorana (solid), $g = 2$ complex scalar (dotted) and $g = 4$ Dirac dark matter (long dash) with DM mass of 10 MeV. The left panel is for $s$-wave annihilation into neutrinos with $\sigma_0 Y_s^{-1/2} = (2.8 - 3.0) \times 10^{-26} \text{ cm}^3/\text{s}$, and the right panel is for $p$-wave annihilation into photons with $\sigma_0 Y_s^{-1/2} = 8.0 \times 10^{-25} \text{ cm}^3/\text{s}$. The horizontal dotted line ($Y_0$) represents the current DM relic density, and $Y_{eq}$ indicates the equilibrium number density.

possibilities (a Dirac fermion, Majorana fermion, complex scalar and real scalar) with a DM mass of 10 MeV. As we expected, the DM number track the equilibrium distribution at very high temperatures, $x_b < 1$. The solution of the Boltzmann equation starts to deviate significantly from the equilibrium abundance at around $x_b \sim 10^{-11}$. Notice that the equilibrium number densities about the temperature are slightly different for particle species. This can be a means to distinguish the nature of the particle.

B. Thermal Equilibrium Approximation

DM particles are assumed to be kept in thermal contact with one of plasmas after neutrino decoupling, and they suddenly decouple at some point. The DM and its products can be expressed by Fermi-Dirac or Bose-Einstein statistics in this case. There can be more particle species in the thermal bath $a$, so we will use the entropy density $s_a \equiv \frac{2}{45} \bar{g}_s T_a^3 = (\rho_a + p_a)/T_a$
to define the number of the effective relativistic degrees of freedom \( \tilde{g}_{*s}^a(T_a) \) in which \( \rho_a \) is the energy density and \( p_a \) is the pressure. The energy density \( \rho_a \) and the pressure \( p_a \) are expressed by

\[
\rho_a = \sum_i \rho_i = \sum_i \frac{g_i}{2\pi^2} \int \frac{dq}{\pi^2} E_i \frac{1}{\exp(E_i/T_a) \pm 1},
\]

\[
p_a = \sum_i p_i = \sum_i \frac{g_i}{2\pi^2} \int dq \frac{q^4}{3E_i} \frac{1}{\exp(E_i/T_a) \pm 1},
\]

where \( g_i \) is the internal degree of freedom for the corresponding particle \( i \), \( E_i = \sqrt{q^2 + m_i^2} \) is the energy with mass \( m_i \) and \( +(-) \) sign is for fermions (bosons). We set the chemical potentials to zero. The number of the effective relativistic degrees of freedom is given by

\[
\tilde{g}_{*s}^a(T_a) = \frac{45}{2\pi^2} \frac{(\rho_a + p_a)}{T_a^4}.
\]

Since the entropy in each thermal bath is conserved after neutrino decoupling, the temperature \( T_a(T_b) \) varies as \( \tilde{g}_{*s}^{a-1/3} R^{-1} \left( \frac{g_{*s}^b}{R_{*s}} \right) \) in which \( R \) is the scale factor. We can find the temperature ratio at the DM decoupling time if we know the temperature ratio at a certain time (the time of neutrino decoupling). The temperature ratio at the DM decoupling time results in

\[
\frac{T_{aD}}{T_{bD}} = \left( \frac{\tilde{g}_{*s}^a (T_{aD})}{\tilde{g}_{*s}^a (T_{bD})} \right)^{1/3} \left( \frac{\tilde{g}_{*s}^b (T_{bD})}{\tilde{g}_{*s}^b (T_{aD})} \right)^{1/3},
\]

where \( T_{aD} \) is the neutrino decoupling temperature in which photon and neutrino temperatures are the same, \( T_{bD} \) is the neutrino or photon temperature \(^4\) at DM decoupling. This formula can be approximated to the temperature ratio at late times, \( i.e., \) temperatures \( T_a, T_b \) less than the decoupling temperatures. Since the DM decoupling occurs at \( x_D \sim 18 \) as we see in the subsection A, there must be almost no DM contribution to the relativistic degrees of freedom \( \tilde{g}_{*s}^a(T_{aD}) \). We remove the tilde. The temperature ratio after the DM decoupling is then given by

\[
\frac{T_a}{T_b} \sim \left( \frac{\tilde{g}_{*s}^a (T_{aD})}{\tilde{g}_{*s}^a (T_a)} \right)^{1/3} \left( \frac{\tilde{g}_{*s}^b (T_{bD})}{\tilde{g}_{*s}^b (T_{aD})} \right)^{1/3}.
\]

We now determine \( N_{\text{eff}} \) in each case. If the DM particle interacts with neutrino \( (a = \nu \) and \( b = \gamma) \), the electromagnetic plasma is not relevant to the DM interaction. we can identify

---

\(^3\) The mark “ \( \sim \) ” places on top of the symbol of the number of the effective relativistic degrees of freedom to indicate the DM inclusion. If there is no mark “ \( \sim \) ”, the DM is excluded.

\(^4\) Notice that one of the DM decoupling temperatures is determined when the equilibrium DM number is the same as the present-day DM relic density, \( Y_{eq}(T_{bD}) = Y_0 \).
FIG. 2: The effective number of neutrino degrees of freedom, $N_{\text{eff}}$, as a function of a thermal dark matter mass $M$. Curves correspond to a $g = 1$ self-conjugate scalar (short dash), $g = 2$ Majorana (solid), $g = 2$ complex scalar (dotted) and $g = 4$ Dirac dark matter (long dash). The upper (lower) curves are for the case when the dark matter particles are in thermal equilibrium with neutrinos (electrons and photons). The dark horizontal band is the Planck CMB 1σ allowed range and the light dark band is the 2σ upper allowed range.

(\frac{g^\gamma_s (T_\gamma)}{g^\nu_s (T_\nu)})^{1/3}$ with $(T_\nu/T_\gamma)^{SM}$. We get the effective number of neutrino species from Eqs. (2) and (9),

$$N_{\nu\text{eff}} = N_{\text{eff}}^{SM} \left( \frac{g^\nu_s (T_\nu)}{g^\nu_s (T_\nu)} \right)^{4/3}. \tag{10}$$

For electromagnetic coupled DM ($a = \gamma$ and $b = \nu$), there is only one species in the neutrino thermal bath, $\nu$. Then, the effective relativistic degrees of freedom $g^\gamma_s (T_\gamma)$ will be the same at any time. Since $(T_\nu/T_\gamma)^{SM} = (g^\gamma_s (T_\gamma)/g^\nu_s (T_\nu))^{1/3}$, we get the effective number of neutrino species

$$N_{\gamma\text{eff}} = N_{\text{eff}}^{SM} \left( \frac{g^\gamma_s (T_\nu)}{g^\nu_s (T_\nu)} \right)^{4/3}. \tag{11}$$

In Fig. 2 we show the numerical result of the $N_{\text{eff}} - M$ relation for several possibilities (a Dirac fermion, Majorana fermion, complex scalar and real scalar) of neutrino coupled DM particles by the upper set of curves and electromagnetically coupled DM particles by
TABLE I: 1σ and 2σ lower limits on the dark matter mass and upper limits on the existence of any other dark radiation for dark matter particles in the thermal equilibrium with neutrinos or electromagnetic plasmas. The mark ‘−’ indicates that the limit is irrelevant. The symbol “S” stands for scalar and “F” for fermion.

| $g$ | Neutrino coupled DM (MeV) | EM coupled DM (MeV) | $\Delta N_{\text{eff}}$ |
|-----|---------------------------|---------------------|------------------|
| 1σ  | 6.4 1 (S)  9.3 2 (S,F)  11.9 4 (F) | 10.1 2 (S)  12.6 2 (S,F)  14.9 4 (F) | 0.94 1 (S)  1.37 2 (S,F)  1.27 2 (F)  1.80 4 (F) |
| 2σ  | 3.7 1 (S)  7.1 2 (S,F)  10.0 4 (F) | – – – – | 1.17 1 (S)  1.60 2 (S,F)  1.50 2 (F)  2.03 4 (F) |

the lower set curves. We have implicitly assumed that the neutrinos decouple at $T_{\nu d} \approx 2.3$ MeV \[25\]~\[27\]. In the case that DM particles are in equilibrium with neutrinos, $N_{\text{eff}}$ increases for lighter DMs. Conversely, $N_{\text{eff}}$ decreases in equilibrium with electromagnetic plasma. We put a bound on the DM mass by requiring that $N_{\text{eff}}$ is compatible with the measured value from Planck \[5\], and it is listed at Table. I. If there is a significant, but small, density of additional radiation, it can be explained by the neutrino coupled DM with mass around 10 MeV. We should notice that the dark radiation predictions in the existence of light DM particles are different for particle species. There is still enough room for the existence of the extra dark radiation particle such as sterile neutrino or Goldstone boson which has its decoupling temperature less than 100 MeV, allowed for the electromagnetically coupled DM.

C. Out-of-Equilibrium Production

As we can see in Fig. 1, there is a smooth transition between two regimes, before and after DM freeze-out, and DM particles does not track significantly the equilibrium from $x_b \sim 10 – 11$. We cannot use the equilibrium thermal dynamics for consideration of the smooth transition. We can instead apply the second law of thermodynamics for the change in entropy of relativistic particles in the adiabatically expanding universe,

\[
dS_a = \frac{dQ}{T_b},
\]

where $dQ = d(R^3 \rho_{\text{DM}})$ is the heat added per comoving volume due to the DM annihilations. Since the number of DM particles is reduced by their annihilation at the temperature less than its mass, the energy density of DM can be described in its nonrelativistic approximation,
\( \rho_{\text{DM}} \simeq n_{\text{DM}} M = M s_b Y \). The change in entropy\(^5\) is given by
\[
dS_a = -S_b x_b dY \rightarrow \Delta S_a = -S_b \int_i x_b dY ,
\]
where \( i \) is an initial point. We consider the initial point at neutrino decoupling time, because it is the last point in which neutrinos and photons are in thermal contact. DM particles must be in thermal equilibrium in the thermal bath \( a \) at the initial point. Our observation point is the recombination time, long time later after the freeze-out. It must be enough time that the produced particles are thermalized in the thermal background. The change in entropy can be then expressed by
\[
\Delta S_a = S_a - S_{ai} = \frac{2\pi^2}{45} \left[ g_{*s}^a T_a^3 R^3 - \left( g_{*s}^a T_a^3 R^3 \right)_i \right] .
\]
(14)

The temperature ratio is determined by a combination of Eqs. (13) and (14),
\[
\left( \frac{T_a}{T_b} \right)^3 = \left( \frac{g_{*s}^a}{g_{*s}^b} \right) \left( \frac{R_i}{R} \right)^3 \left( \frac{T_{ai}}{T_b} \right)^3 - \frac{g_{*s}^b}{g_{*s}^a} \int_i x_b dY ,
\]
(15)
where \( T_{ai} \) is, according to Ref. [6], almost the same as the neutrino decoupling temperature described in the SM of the DM absence. Using the entropy conservation \( (g_{*s}^a R^3 \sim T^{-3}) \) in the SM, we can approximate the first term of Eq. (15). Then the temperature ratio is given by
\[
\left( \frac{T_a}{T_b} \right)^3 \simeq \left( \frac{T_a}{T_b} \right)^3_{SM} - \frac{g_{*s}^b}{g_{*s}^a} \left[ x_b Y - (x_b Y)_i - \int_i Y dX \right] ,
\]
(16)
where we introduced the integration method by parts in convenience for numerical computations. The first term on the right-hand side of Eq. (16) is just the original temperature ratio in radiation, and the second term represents a contribution by the DM annihilation. We can then express the temperature ratio in each case, neutrino coupled DM \((a = \nu \text{ and } b = \gamma)\) and electromagnetic coupled DM \((a = \gamma \text{ and } b = \nu)\).

We show the numerical result of the \( N_{\text{eff}} - M \) relation for neutrino coupled DM by the upper set of curves and electromagnetically coupled DM by the lower set curves in Fig. 3. The basic arguments are the same as in the equilibrium approximation of the subsection B. The bounds on the DM mass with the measured value from Planck are also listed at Table II, as well as possibility of the existence of extra dark radiation particles for the electromagnetically coupled DM. This method gives much tighter constraint to the bounds of the DM

\(^5\) Notice that the plasma in the thermal bath \( b \) is always in the thermal equilibrium because it is not relevant to the DM interaction, and so the entropy \( S_b \) is constant.
FIG. 3: Same as Fig. 2, but contour lines for the case when radiation particles are produced in nonequilibrium method (freeze-out mechanism).

TABLE II: Same as Table. I, but the values for the case when radiation particles are produced in nonequilibrium method (freeze-out mechanism).

| $g$ | Neutrino coupled DM (MeV) | EM coupled DM (MeV) | $\Delta N_{\text{eff}}$ |
|-----|--------------------------|-------------------|-----------------|
|     | 1 (S) 2 (S,F) 4 (F)     | 1 (S) 2 (S,F) 4 (F) | 1 (S) 2 (S) 2 (F) 4 (F) |
| 1σ  | 9.6 12.3 14.8           | 9.1 11.8 14.3     | 0.70 0.99 0.94 1.36 |
| 2σ  | 7.3 10.3 12.9           | – – –             | 0.93 1.22 1.16 1.56 |

mass, and also gives tighter constraint to the existence of additional dark radiation particles for electromagnetically coupled DM. We interpret the reason in the following way. DM particles annihilate more slowly and smoothly into SM particles. The slower annihilation results in the smaller expansion, eventually smaller size of the universe later. The same amount of relativistic particle number must be produced by DM annihilations in equilibrium and nonequilibrium methods. The predictions of energy density are however different, because their annihilating period is different.
III. CONCLUSIONS

The light ($M \leq 20$ MeV) dark matter particles freeze out after neutrino decoupling. If dark matter particles interact with neutrino or electromagnetic plasma, the late time entropy production can change the neutrino-to-photon temperature ratio, equally effective number of neutrinos $N_{\text{eff}}$. If there is a significant, but small, density of additional radiation, it can be explained by the neutrino coupled DM with mass around 10 MeV. The effective number of neutrino species $N_{\text{eff}}$ is reduced by photon heating. In this case, the existence of additional dark radiation particles can help to interpret agreement of the current observations. The dark matter particles are nonrelativistic when they decouple, so we used the freeze-out (out-of-equilibrium) mechanism besides the thermal equilibrium approximation. The dark matter particles smoothly annihilate into the SM particles. The slower annihilation eventually results in the smaller expansion rate (eventually smaller size of the universe later). Although the same amount of relativistic particle number is produced by dark matter annihilations in the equilibrium approximation and the nonequilibrium method, the predictions of the energy density at late times are different. We demonstrated that the bound of dark matter mass and the possibility of the existence of extra radiation particles were more tightly constrained in the nonequilibrium method.

Acknowledgments

The work is supported by the National Research Foundation of Korea (NRF) grant funded by Korea government of the Ministry of Education, Science and Technology (MEST) (Grant No. 2011-0017430) and (Grant No. 2011-0020333).
[1] D.A. Dicus, E.W. Kolb, A.M. Gleeson, E.C.G. Sudarshan, V.L. Teplitz, M.S. Turner, Phys. Rev. D 26, 2694 (1982).

[2] G. Mangano, G. Miele, S. Pastor, M. Peloso, Phys. Lett. B 534, 8 (2002), astro-ph/0111408.

[3] K.N. Abazajian, M.A. Acero, S.K. Agarwalla, et al., arXiv:1204.5379.

[4] S. Weinberg, Phys. Rev. Lett. 110, 241301 (2013), arXiv:1305.1971.

[5] P.A.R. Ade, et al. [Planck Collaboration], arXiv:1502.01589.

[6] C.M. Ho, R.J. Scherrer, Phys. Rev. D 87, 023505 (2013), arXiv:1208.4347; Phys. Rev. D 87, 065016 (2013), arXiv:1212.1689.

[7] G. Steigman, Phys. Rev. D 87, 103517 (2013), arXiv:1303.0049.

[8] C. Boehm, M.J. Dolan, C. McCabe, J. Cosmol. Astropart. Phys. 08 (2013) 041, arXiv:1303.6270.

[9] K.M. Nollett, G. Steigman, Phys. Rev. D 89, 083508 (2014), arXiv:1312.5715; arXiv:1411.6005.

[10] A.G. Riess, L. Macri, S. Casertano, H. Lampeitl, H.C. Ferguson, A.V. Filippenko, S.W. Jha, W. Li, et al., Astrophys. J. 730, 119 (2011) [Erratum-ibid. 732, 129 (2011)], arXiv:1103.2976.

[11] J. Yang, D. Schramm, G. Steigman, R.T. Rood, Astrophys. J. 227, 697 (1979).

[12] R. Cooke, M. Pettini, R.A. Jorgenson, M.T. Murphy, C.C. Steidel, Astrophys. J. 781, 31 (2014), arXiv:1308.3240.

[13] M. Pettini, R. Cooke, Mon. Not. R. Astron. Soc. 425, 2477 (2012), arXiv:1205.3785.

[14] E. Aver, K.A. Porter, R.L. Porter, E.D. Skillman, J. Cosmol. Astropart. Phys. 11, 017 (2013), arXiv:1309.0047.

[15] Y.I. Izotov, T.X. Thuan, N.G. Guseva, Mon. Not. Roy. Astron. Soc. 445, 778 (2014), arXiv:1408.6953.

[16] E.W. Kolb, M.S. Turner, T.P. Walker, Phys. Rev. D 34, 2197 (1986).

[17] P.D. Serpico, G.G. Raffelt, Phys. Rev. D 70, 043526 (2004), astro-ph/0403417.

[18] C. Boehm, P. Fayet, Nucl. Phys. B 683, 219 (2004), hep-ph/0305261.

[19] C. Boehm, D. Hooper, J. Silk, M. Casse, Phys. Rev. Lett. 92, 101301 (2004), astro-ph/0309686.

[20] D. Hooper, F. Ferrer, C. Boehm, J. Silk, J. Paul, N.W. Evans, M. Casse, Phys. Rev. Lett. 93,
161302 (2004), astro-ph/0311150.

[21] K. Ahn, E. Komatsu, Phys. Rev. D 72, 061301 (2004), astro-ph/0506520.

[22] C. Boehm, M.J. Dolan, C. McCabe, J. Cosmol. Astropart. Phys. 12, 027 (2012), arXiv:1207.0497.

[23] D.P. Finkbeiner, S. Galli, T. Lin, T.R. Slatyer, Phys. Rev. D 85, 043522 (2012), arXiv:1109.6322.

[24] L. Lopez-Honorez, O. Mena, S. Palomares-Ruiz, A.C. Vincent, J. Cosmol. Astropart. Phys. 07, 046 (2013), arXiv:1303.5094.

[25] K. Enqvist, K. Kainulainen, V. Semikoz, Nucl. Phys. B 374, 392 (1992).

[26] A.D. Dolgov, Phys. Rept. 370, 333 (2002), hep-ph/0202122.

[27] S. Hannestad, Phys. Rev. D 65, 083006 (2002), astro-ph/0111423.