We review the impacts of the recent precision cosmic microwave background anisotropy measurements on particle physics. Topics include constraints on the axion, dark radiation, and annihilating dark matter.

Subject Index E63, E70, E74, E75

1. Introduction

After the first measurements of the cosmic microwave background (CMB) anisotropy by the COBE satellite [1], there has been great progress with CMB observations. The WMAP satellite took a very important role in using the CMB as a tool for exploring high energy physics [2]. Recently, the Planck satellite provided us with a detailed map of the CMB with unprecedented accuracy [3]. Currently, many ground-based CMB detectors, especially dedicated for the detection of CMB polarization, are in action.

Precision measurement of the CMB has made it possible to probe/constrain high energy physics. Currently, inflation models are severely restricted. This is far beyond the capability of accelerator experiments and it shows the power of CMB as a probe into high energy physics. There are also many implications of the precise CMB measurements for particle physics. In this article, we provide a brief overview of the impacts of recent precise CMB observations on the axion, dark radiation, and annihilating dark matter (DM).

2. Axion

The axion was introduced as a solution to the strong CP problem in quantum chromodynamics (QCD) [4,5]. In QCD, there exists the following $\theta$-term in the Lagrangian,

$$\mathcal{L} = \frac{\theta g_s^2}{32\pi^2} G^a_{\mu\nu} F^{\mu\nu} a^a,$$

where $g_s$ is the QCD gauge coupling constant, $G^a_{\mu\nu}$ is the field strength of the gluon, and $\theta$ is a constant. This contributes to the neutron electric dipole moment and it gives a constraint on $\theta$ of $\theta \lesssim 10^{-10}$ [6], despite the fact that $\theta$ is naturally expected to be $O(1)$. This is the strong CP problem.
In the presence of the axion, it is replaced with the dynamical term
\[ \theta \rightarrow \theta_{\text{eff}} \equiv \theta - \frac{a}{f_a}, \] (2)
where \( a \) represents the axion and \( f_a \) is the axion decay constant. By minimizing the scalar potential of the axion, it is found that the effective theta angle vanishes: \( \theta_{\text{eff}} = 0 \), hence it solves the strong CP problem.

In order for this mechanism to work, the axion must obtain its mass only through the QCD instanton effect. Otherwise, the minimum of the axion potential deviates from the point at \( \theta_{\text{eff}} = 0 \). To control the mass term of the axion, it is assumed that the axion arises as a pseudo-Nambu–Goldstone boson from the spontaneous breaking of the Peccei–Quinn (PQ) symmetry. The axion mass is given by
\[ m_a \simeq 6 \times 10^{-6} \text{eV} \left( \frac{10^{12} \text{GeV}}{f_a} \right). \] (3)
Phenomenologically, the PQ breaking scale, or the axion decay constant, \( f_a \), is constrained to \( 10^9 \text{GeV} \lesssim f_a \lesssim 10^{12} \text{GeV} \) [7–9].

The axion not only solves the strong CP problem, but can also explain the observed abundance of DM in the universe. The axion begins a coherent oscillation around the QCD phase transition, and its abundance in terms of the density parameter \( \Omega_a \) is estimated as [10]
\[ \Omega_a h^2 \simeq 0.2 \theta_i^2 \left( \frac{f_a}{10^{12} \text{GeV}} \right)^{1.19}, \] (4)
where \( \theta_i \) denotes the initial misalignment angle of the axion, which is the initial condition set during inflation, if PQ symmetry is already broken during inflation.\(^1\) Thus it can explain the observed DM abundance for \( f_a \sim 10^{12} \text{GeV} \) and \( \theta_i \sim O(1) \). Much larger \( f_a \) is allowed if \( \theta_i \ll 1 \) or if late-time entropy production occurs [11].

One of the characteristic properties of the axion DM is that it obtains isocurvature perturbations during inflation, which eventually become the cold DM (CDM) isocurvature perturbation [12–14]. However, the magnitude of the CDM isocurvature perturbation is constrained from CMB observations. Let us discuss this below.

### 2.1. Axion isocurvature perturbation

Since the axion is massless during inflation, it obtains quantum fluctuations at superhorizon scales as
\[ \langle \delta a_k \delta a_{k'} \rangle = (2\pi)^3 \delta^{(3)}(\vec{k} + \vec{k}') \frac{H_{\text{inf}}^2}{2k^3}, \] (5)
where \( H_{\text{inf}} \) is the Hubble scale during inflation. Taking account of the fluctuation, the axion abundance should be evaluated as [15]
\[ \Omega_a h^2 \simeq 0.2 \left[ \theta_i^2 + \left( \frac{H_{\text{inf}}}{2\pi f_a} \right)^2 \right] \left( \frac{f_a}{10^{12} \text{GeV}} \right)^{1.19}. \] (6)

The non-linear axion isocurvature perturbation is defined as [16,17]
\[ S_a \equiv 3(\xi_a - \xi), \] (7)
where \( \xi_a \) and \( \xi \) are the curvature perturbations on the constant axion energy density surface and constant total energy density surface (uniform density slice), respectively.

\(^1\) Hereafter we only consider this case. It is also possible that PQ symmetry is broken after inflation. See the review article [9] for detailed constraints on these cosmological scenarios.
The power spectrum of the CDM isocurvature perturbation is given by
\[
\langle S_C(k) S_C(k') \rangle = (2\pi)^3 \delta^{(3)}(\mathbf{k} + \mathbf{k}') P_{S_C}(k),
\]
where the power spectrum \( P_{S_C}(k) \) is given by
\[
P_{S_C}(k) = \frac{k^3}{2\pi^2} P_{S_C}(k) = 4 r^2 \frac{(H_{\text{inf}}/2\pi)^2}{(\theta_i/\Omega_0^a)^2 + (H_{\text{inf}}/2\pi)^2}.
\]

Here \( r \) is the ratio between the axion energy density and the CDM energy density: \( r = \Omega_{a}^\Omega / \Omega_{C}^\Omega \). According to the WMAP7 results, the bound on the CDM isocurvature perturbation reads \( \alpha \equiv P_{S_C}/P_{\zeta} < 0.15 \) at the reference scale \( k_* = 0.002 \) Mpc\(^{-1} \), where \( P_{\zeta} = 2.43 \times 10^{-9} \) is the dimensionless power spectrum of the curvature perturbation [18]. This severely restricts the inflation energy scale. Figure 1 shows the constraints on the \( (H_{\text{inf}}, \theta_i) \) plane for \( f_a = 10^{11} \) GeV (top) and \( f_a = 10^{15} \) GeV (bottom) [19]. For the standard value of \( f_a \sim 10^{12} \) GeV and \( \theta_i \sim \mathcal{O}(1) \), the inflation scale must be very low: \( H_{\text{inf}} \lesssim 10^7 \) GeV. This gives strong constraints on the inflation model-building if the axion is present in the universe.

2.2. Non-Gaussian axion isocurvature perturbation

The axion isocurvature perturbation also produces non-Gaussian perturbations in the CDM isocurvature perturbation since the axion energy density contains a term like \( (\delta a)^2 \) [16,17,20]. (See also
The non-Gaussian nature of the perturbation is imprinted in the bispectrum as
\[
\langle S_C(k_1)S_C(k_2)S_C(k_3) \rangle = (2\pi)^3 \delta^{(3)}(\vec{k}_1 + \vec{k}_2 + \vec{k}_3) B_{SC}(k_1, k_2, k_3),
\]
(10)
where
\[
B_{SC}(k_1, k_2, k_3) = f_{SC}[P_{SC}(k_1)P_{SC}(k_2) + \text{perms.}].
\]
(11)
As shown in Ref. [16], this gives a different pattern in the CMB temperature fluctuation from that induced by the non-Gaussian adiabatic perturbation.

The constraint using the WMAP seven-year data reads [19]
\[
|f_{NL}^{SS}| = |\alpha f_s| < 140.
\]
(12)
This constraint is also shown in Fig. 1. We find that the bispectrum gives a comparable constraint on the axion model parameters to that from the power spectrum for small $\theta_i$.

3. Dark radiation

Dark radiation consists of relativistic particles which do not effectively interact with photons or baryons except for standard model neutrinos. The total relativistic energy density $\rho_r$ is parameterized by the so-called effective number of relativistic degrees of freedom $N_{\text{eff}}$ as
\[
\rho_r = \left[1 + N_{\text{eff}} \frac{7}{8} \left(\frac{4}{11}\right)^{4/3}\right] \rho_{\gamma},
\]
(13)
where $\rho_{\gamma} = (\pi^2/15)T_{\gamma}^4$ is the photon energy density after the $e^+e^-$ annihilation, with $T_{\gamma}$ representing the photon temperature. The standard model predicts $N_{\text{eff}} = 3.046$ [25]. Thus deviation of $N_{\text{eff}}$ from 3 would imply the existence of light particles beyond the standard model. In the following we denote the dark radiation particle by $X$.

Increasing $N_{\text{eff}}$ leads to several effects on the CMB anisotropy. First, the early integrated Sachs–Wolfe effect around the matter–radiation equality becomes efficient and the first peak is enhanced. Second, since the sound horizon at the recombination epoch becomes small, or $H_{\text{rec}}$ becomes large, the peak position in the CMB power spectrum shifts to the small scale. The anisotropic stress effect and the enhanced Silk damping due to the increased Hubble expansion also cause non-trivial effects on the CMB anisotropy [2]. The Planck satellite result combined with the CMB polarization measurement by WMAP and small-scale CMB measurement by the SPT and ACT give $N_{\text{eff}} = 3.36^{+0.68}_{-0.64}$ [3].

On the other hand, $N_{\text{eff}}$ is also constrained from the observation of primordial light element abundance. Since increasing $N_{\text{eff}}$ makes the Hubble expansion rate larger at the freezeout of $p \leftrightarrow n$ conversion, the $n/p$ ratio becomes larger, leading to larger primordial helium abundance. Recent analysis of the H II regions shows a slight preference for the non-standard value $N_{\text{eff}} = 3.51 \pm 0.35$ [26].

3.1. Models of dark radiation

Dark radiation is predicted in many theories beyond the standard model. Here we present some explicit models of dark radiation.

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Note that the epoch of matter–radiation equality is delayed, or $H_{\text{eq}}$ becomes smaller, and hence the peak position of the matter power spectrum shifts to the large scale.
Table 1. Charge assignments for fermions $\Psi$ and bosons $\Phi$.

| $SO(10) \times U(1)_\psi$ | $\Psi_{16}(1)$ | $\Psi_{10}(-2)$ | $\Psi_1(4)$ |
| $SU(5) \times U(1)_\psi \times U(1)_x$ | $\psi_{10}(1, 1)$ | $\psi_5(1, -3)$ | $\psi_{10}(-2)$ | $\psi_1(4, 0)$ |
| $SO(10) \times U(1)_\psi$ | $\Phi_{16}(1)$ | $\Phi_{10}(-2)$ | $\Phi_1(4)$ |
| $SU(5) \times U(1)_\psi \times U(1)_x$ | $\phi_{10}(1, 1)$ | $\phi_5(1, -3)$ | $\phi_{10}(-2)$ | $\phi_1(4, 0)$ |

3.1.1. “Thermal” dark radiation. First, let us consider the dark radiation produced thermally, i.e., $X$ was in thermal equilibrium in the early universe and at some epoch it decouples from the thermal bath [27–31]. In this case, the extra relativistic degrees of freedom are given by

$$\Delta N_{\text{eff}} = \frac{4}{7} \left[ \frac{10.75}{g_{\text{eff}}(T_{\text{dec}})} \right]^{4/3} \times \begin{cases} 1 & \text{for a real scalar} \\ 7/4 & \text{for a Weyl fermion} \\ 2 & \text{for a gauge boson} \end{cases}$$

(14)

where $T_{\text{dec}}$ is the decoupling temperature of $X$ and $g_{\text{eff}}$ is the relativistic degrees of freedom at that temperature. We have $\Delta N_{\text{eff}} \sim 1$ if the decoupling temperature is around 10–100 MeV. In Ref. [28] it is shown that a real scalar and a gauge boson are not favored as a candidate of $X$ to explain $\Delta N_{\text{eff}} \sim 1$ by taking account of the constraints from the arguments of supernova cooling.

Let us suppose that $X$ is a chiral fermion and it has an effective interaction with the standard fermions $f$ as

$$\mathcal{L}_{\text{eff}} = \frac{1}{\Lambda^2} (\bar{f} \gamma^\mu f) (\bar{X} \gamma_\mu X),$$

(15)

where $\Lambda$ is the cutoff scale. By considering the interaction of $X$ with electrons, the decoupling temperature may be estimated as

$$T_{\text{dec}} \sim 10 \text{ MeV} \left( \frac{\Lambda}{1 \text{ TeV}} \right)^{4/3}.$$ 

(16)

Therefore, a cutoff scale of the order of TeV scale will be needed. It can marginally satisfy the bound from star cooling and LHC.

An example is the $E_6$-motivated extension of the standard model [28]. We introduce a new $U(1)$ gauge symmetry, under which $X$ is charged. It forbids the mass term of $X$ and ensures the smallness of its mass. The anomaly-free $U(1)$ other than the B-L is obtained by considering the $E_6$ gauge group as a starting point. Let us consider the symmetry-breaking pattern $E_6 \rightarrow SO(10) \times U(1)_\psi \rightarrow SU(5) \times U(1)_\psi \times U(1)_x$. The field contents and their charge assignments are summarized in Table 1. Both the fermions $\Psi$ and bosons $\Phi$ are contained in the fundamental representation 27 of $E_6$. Standard model fermions are included in $\psi_{10}$ and $\psi_5$, and right-handed neutrinos correspond to $\psi_1$. Standard model Higgs bosons are contained in $\phi_5$ and $\phi_8$. The dark radiation candidate is a singlet fermion $\psi_1$ which is charged under only extra $U(1)$. The allowed term is

$$\mathcal{L} = \frac{1}{M} \bar{\phi}_1^* \phi_1^* \psi_1 \psi_1 + \text{h.c.},$$

(17)

which becomes the mass term after $\phi_1$ obtains a VEV. If $\langle \phi_1 \rangle \sim O(1)$ TeV, we obtain $\Lambda \sim O(1)$ TeV and $m_X \sim 10^{-3}$ eV for $M = M_F$. Therefore it is a good candidate for dark radiation. In this model a new gauge boson, which we call $Z'$, appears at the mass of $O(1)$ TeV. The non-detection of the $Z'$-boson at the LHC will ensure the exclusion of this scenario.

Another scenario is the so-called self-interacting dark radiation based on the unbroken hidden gauge symmetry [31]. Let us introduce a standard model singlet scalar $\phi$ charged under only the
hidden U(1)' does not have kinetic mixing with the standard model U(1) as long as the U(1)' is not spontaneously broken. It can have a renormalizable coupling to the standard model Higgs boson $H$ through

$$\mathcal{L} = \frac{\lambda}{4} |\phi|^2 |H|^2.$$  

(18)

Integrating out $\phi$, we obtain the effective Lagrangian

$$\mathcal{L}_{\text{eff}} = \frac{1}{\Lambda_1^2} F'_{\mu\nu} F'_{\mu\nu} \frac{|H|^2}{\Lambda_1},$$  

(19)

where $F'_{\mu\nu}$ represents the field strength of the hidden U(1)' and $\Lambda_1 \sim (\lambda g' G/8\pi^2)^{-1/2} m_\phi$, with $g'$ being the hidden U(1)' gauge coupling constant and $m_\phi$ being the mass of $\phi$. In order to avoid too large a Higgs invisible decay width into the hidden gauge boson, we need $\Lambda_1 \phi \sim (\lambda g' G/8\pi^2)^{-1/2} m_\phi$, with $g'$ being the hidden U(1)' gauge coupling constant and $m_\phi$ being the mass of $\phi$. In order to avoid too large a Higgs invisible decay width into the hidden gauge boson, we need $\Lambda_1 \phi \sim (\lambda g' G/8\pi^2)^{-1/2} m_\phi$. Hidden gauge bosons can be thermalized through this effective interaction for $T < m_\phi$, and the decoupling temperature is estimated as

$$T_{\text{dec}} \sim 200 \text{ GeV} \left(\frac{100}{g_*(T_{\text{dec}})}\right)^{1/6} \left(\frac{\Lambda_1}{10^6 \text{ GeV}}\right)^{4/3}.$$  

(20)

In the case of hidden U(1)' with no additional chiral matter, we obtain at most $\Delta N_{\text{eff}} \sim 0.1$. By extending the hidden gauge group to non-Abelian ones, or by adding extra chiral fermions charged under the hidden gauge group, we can have larger $\Delta N_{\text{eff}}$. In such cases, the dark radiation sector has self interactions and hence has smaller viscosity [32] compared with free-streaming relativistic particles, which may leave distinct signatures in the CMB.

### 3.1.2. “Nonthermal” dark radiation

Next let us consider the case where a heavy particle $\phi$ decays into dark radiation $X$ [33–52]. The extra relativistic degrees of freedom are given by

$$\Delta N_{\text{eff}} \simeq 43 \left[\frac{10.75}{g_*(T_{\phi})}\right]^{1/3} \left[\frac{B_X \rho_\phi}{\rho_\gamma}\right]_{H=\Gamma_\phi}.$$  

(21)

where $T_{\phi}$ is the temperature at the $\phi$ decay, $\Gamma_\phi$ is the decay rate of $\phi$, and $B_X$ denotes the branching fraction of $\phi$ into $X$ particles. Here we have assumed that $\phi$ decays in the radiation-dominated era. Thus if $\phi$ decays when it (nearly) dominates the universe, we may obtain $\Delta N_{\text{eff}} \sim 1$.

A well-motivated example is the saxion in supersymmetric axion models as $\phi$, which decays into a pair of axions as $X$ [35]. The decay rate is given by

$$\Gamma_{\phi \rightarrow 2a} \simeq \frac{1}{64 \pi} \frac{m_\phi^3}{f_a^2}.$$  

(22)

This is often the dominant decay mode of the saxion since the decay into gauge bosons is suppressed by the loop factor. Assuming that the saxion begins a coherent oscillation around $H = m_\phi$ with an initial amplitude $\phi_i$, we obtain

$$\left[\frac{\rho_\phi}{\rho_\gamma}\right]_{H=\Gamma_\phi} \simeq \frac{T_R}{6T_\phi} \left(\frac{\phi_i}{M_P}\right)^2 \simeq 0.6 \left(\frac{T_R}{10^7 \text{ GeV}}\right) \left(\frac{10^3 \text{ TeV}}{m_\phi}\right)^{3/2} \left(\frac{f_a}{10^{10} \text{ GeV}}\right) \left(\frac{\phi_i}{M_P}\right)^2,$$  

(23)

where $T_R$ denotes the reheating temperature of the universe. Thus the axions produced by the saxion decay can be dark radiation for an appropriate choice of parameters [35].

[3] The saxion dynamics may be significantly changed by thermal effects [39,40,50,53].
Recently, an interesting possibility has been proposed that the string moduli decays into the axion pair, which form the dark radiation $X$ [54–57]. Let us consider the type IIB string theory and suppose that the moduli stabilization after the compactification is achieved in a shift symmetric way:

$$T \rightarrow T + i\alpha,$$

(24)

where $T$ is the closed string modulus superfield with $\alpha$ being a real parameter. This can be regarded as a remnant of higher-dimensional gauge symmetry for the Ramond–Ramond field. Then the four-dimensional theory should contain the Kähler potential of the form $K(T + T^\dagger)$. In this setup, the imaginary component of the modulus, which we call axion, remains massless. The modulus decays into its axionic partner with the decay rate [56]

$$\Gamma_{\tau \rightarrow 2a} = \frac{1}{64\pi} \frac{K_{TTT}^2}{K_{TT}^3} m_\tau^3,$$

(25)

where we have defined the canonically normalized moduli as $T - \langle T \rangle \equiv (\tau + i\alpha)/\sqrt{2K_{TT}}$. Often this is the dominant decay channel of the modulus as shown in Refs. [54–57] in some explicit moduli stabilization scenarios. In general, this causes the so-called “moduli-induced axion problem” [56] since the moduli tend to dominate the universe before the decay. One way to avoid the axion overproduction is to introduce the Giudice–Masiero term in the Kähler potential:

$$K = -3\log \left[ T + T^\dagger - \frac{1}{3} \left( |H_u|^2 + |H_d|^2 + (z H_u H_d + \text{h.c.}) \right) \right],$$

(26)

with $z$ being a numerical constant. This leads to the modulus decay into the Higgs boson pair, and then the branching fraction in to the axion pair is evaluated as $B_X = 1/(1 + 2z^2)$. Thus by choosing the value of $z$, we can obtain $\Delta N_{\text{eff}} \sim 1$. It may be surprising that the cosmological observations can give a constraint on the string model-building, in particular the model of compactification. Figure 2 shows the constraint on the modulus decay rate into the axion pair $\Gamma_a$ and that into the standard model particles $\Gamma_{SM}$ [56]. Further ideas to detect or constrain the axionic dark radiation have been proposed [58–60].

**Fig. 2.** Constraint on the modulus decay rate into the axion pair $\Gamma_a$ and that into the standard model particles $\Gamma_{SM}$ [56].
3.2. **Dark radiation isocurvature perturbation**

Depending on the production mechanism, the dark radiation $X$ can have an isocurvature perturbation [43,48]. For example, in the nonthermal scenario considered above, the axion dark radiation has an isocurvature perturbation if the parent saxion field obtains quantum fluctuations during inflation—see also Ref. [45] for a model predicting a dark radiation isocurvature perturbation. The dark radiation isocurvature perturbation behaves as the neutrino density isocurvature mode and its magnitude is constrained from observations.

We define the power spectrum of the dark radiation isocurvature perturbation as

$$
\langle S_{\text{DR}}(k) S_{\text{DR}}(k') \rangle = (2\pi)^3 \delta^{(3)}(\vec{k} + \vec{k}') P_{\text{SDR}}(k).
$$

(27)

Defining $\alpha_{\text{DR}} \equiv P_{\text{SDR}}/P_\zeta$, we obtain the constraint $\alpha_{\text{DR}} < 0.16$ at the 95% C.L. from the combination of the WMAP seven-year data and the ACT, if there is no correlation between the isocurvature and adiabatic modes [43].

Moreover, the dark radiation isocurvature perturbation generally has a non-Gaussian component and it can be enhanced in some cases. Future constraints on the non-Gaussian dark radiation isocurvature perturbation are studied in Ref. [48].

4. **Dark matter annihilation**

In the standard weakly interacting massive particle (WIMP) scenario, dark matter particles were in thermal equilibrium in the early universe, and then freeze out due to the expansion of the universe at the temperature $T_{\text{fr}} \sim m_{\text{DM}}/20$. The relic DM abundance is determined by its annihilation cross section as

$$
\Omega_{\text{DM}} h^2 \simeq 0.1 \left( \frac{3 \times 10^{-26} \text{cm}^3 \text{s}^{-1}}{\langle \sigma v \rangle} \right).
$$

(28)

If there is a late-decaying particle such as the gravitino and moduli, which we collectively express by $\phi$, the DM abundance is given by

$$
\Omega_{\text{DM}} h^2 \simeq \min \left[ \frac{m_{\text{DM}}}{m_\phi} \Omega_\phi h^2, \begin{array}{c} 0.1 \left( \frac{3 \times 10^{-26} \text{cm}^3 \text{s}^{-1}}{\langle \sigma v \rangle} \right) \left( \frac{T_{\text{fr}}}{T_\phi} \right) \end{array} \right],
$$

(29)

where $T_\phi$ is the decay temperature of $\phi$ and $\Omega_\phi$ is the abundance of $\phi$. In such a case, the observed DM abundance is explained for larger annihilation cross section.

The DM annihilation cross section is constrained from various observations. Cosmic rays such as gamma-rays, positrons, anti-protons, and neutrinos are produced by the DM annihilation in our Galaxy or other galaxies. If the DM annihilation yields hadrons, gamma-ray and anti-proton observations give very tight constraints.

The DM annihilation cross section is also constrained cosmologically. The success of big-bang nucleosynthesis (BBN) constrains energy injection during the BBN epoch [61–63]. In the case of hadronic annihilation, BBN gives relatively strong constraints. On the other hand, if the DM annihilates dominantly into leptons, the BBN constraint is weak. Even in such a case, the CMB observation gives a stringent constraint. This is because the photons and electrons emitted by DM annihilation ionize neutral hydrogen formed around the epoch of recombination and hence broaden the “width” of the last scattering surface which suppresses the CMB temperature anisotropy at high multipoles and also increases the power of CMB polarization at low multipoles [64,65]. The evolution of the
Fig. 3. CMB temperature power spectrum with no DM annihilation (solid), DM annihilation with \(\langle \sigma v \rangle = 10^{-24} \text{ cm}^3 \text{ s}^{-1}\) (dotted), and \(\langle \sigma v \rangle = 10^{-23} \text{ cm}^3 \text{ s}^{-1}\) (dashed) [68]. It is assumed that DM annihilates into the \(e^+ e^-\) pair.

Fig. 4. Constraint on the DM annihilation cross section in the case of DM annihilation into \(e^+ e^-\) (top) and \(W^+ W^-\) (bottom) [68].

ionization fraction of the hydrogen, \(x_e\), includes the additional term

\[
- \left[ \frac{d x_e}{d z} \right]_{\text{DM}} = \int \frac{dz'}{H(z')(1+z')} \frac{n_{\text{DM}}(z')(\sigma v)}{n_H(z')} \frac{m_{\text{DM}}}{E_{\text{Ry}}} \frac{d \chi_{\text{ion}}^{(F)}(m_{\text{DM}}, z', z)}{dz},
\]  

(30)
with
\[
\frac{d\chi_{\text{ion}}^{(F)}(m_{\text{DM}}, z', z)}{dz} = \int dE \frac{E}{m_{\text{DM}}} \left[ \frac{dN_e^{(F)}}{dE} \frac{d\chi_{\text{ion}}^{(e)}(E, z', z)}{dz} + \frac{1}{2} \frac{dN_{\gamma}^{(F)}}{dE} \frac{d\chi_{\text{ion}}^{(\gamma)}(E, z', z)}{dz} \right],
\]
(31)
where \(n_H\) is the hydrogen number density, \(E_{\text{Ry}} = 13.6 \text{ eV}\) is the Rydberg energy, \(dN_e^{(F)}/dE\) denotes the spectrum of electrons and photons produced by the cascade decay of \(F(= W^+ + W^-, \text{ etc.})\), and \(d\chi_{\text{ion}}^{(i)}(E, z', z)\) is the energy fraction of the injected electron \((i = e^\pm)\) and photon \((i = \gamma)\) energy \(E\) at the redshift \(z'\) which is used for the ionization of hydrogen at the redshift between \(z\) and \(z + dz\).

This was calculated in detail in Refs. [66,67].

Figure 3 shows the temperature power spectrum with and without DM annihilation effects [68]. Thus, precise observations of the CMB anisotropy constrain the DM annihilation rate around the recombination epoch [65–75].

Moreover, the annihilation cross section may have non-trivial velocity dependence if the annihilation takes place through the S-channel resonance or the Sommerfeld enhancement effect. Phenomenologically we parametrize the velocity-dependent annihilation cross section as
\[
\langle \sigma v \rangle = \langle \sigma v \rangle_0 \epsilon + (v/v_0)^n,
\]
(32)
with \(v_0 = 0.3\). Figure 4 shows a constraint on the DM annihilation cross section using the WMAP five-year data in the case of DM annihilation into \(e^+e^-\) (top) and \(W^+W^-\) (bottom) [68] for \(n = 0\) (velocity-independent) and \(n = 1\) with various values of \(\epsilon\).

5. Conclusions
In this article we have reviewed some implications of the recent precise measurements of CMB anisotropy for particle physics. In particular, we focused on constraints on the axion physics, dark radiation, and annihilating dark matter. This shows that CMB measurement is a powerful tool for exploring the physics of the dark sector. Actually, it can even probe the dark sector that accelerator experiments cannot reach.

The physics of CMB is so rich that there are many other applications to high energy physics which have been omitted from this article. In particular, inflation models are constrained from precise measurements of CMB. The non-Gaussianity in the temperature fluctuation is also useful to constrain the model to create the curvature perturbation. These topics are covered by other review articles.

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References


