

Phase transitions in the early Universe: Impacts on BBN and CMB observables

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(Received 1 August 2025; accepted 2 November 2025; published 24 November 2025)

We explore the effects of low-scale cosmological first-order phase transitions (PTs) on big bang nucleosynthesis (BBN) and cosmic microwave background (CMB) anisotropies and distortions. We examine two scenarios: the PT energy acts as additional background energy, modifying the expansion history; and serves as energy injection into the baryon-photon fluid, generating CMB spectral distortions while having a negligible effect on the expansion history. Our analysis reveals that the CMB spectrum from the Planck 2018 dataset imposes stringent constraints on sub-MeV-scale PTs, while sub-keV-scale PTs can induce significant CMB distortions. Consequently, BBN and CMB observables serve as complementary tools to constrain PT parameters, alongside gravitational wave detections.

DOI: [10.1103/s5kg-lhw8](https://doi.org/10.1103/s5kg-lhw8)

I. INTRODUCTION

Cosmological phase transitions (PTs) have quite a few effects on the evolution of the universe, and these effects are basically produced by first-order PTs [1], like baryogenesis [2], magnetic fields [3,4]. First-order PTs also have the potential to generate gravitational waves (GWs). The release of latent heat and associated bubble collisions or turbulence during PT can produce a stochastic background of primordial GWs [5–13]. These GWs carry information about early Universe physics beyond the Standard Model [14–16], and could be detected by LISA [17,18], Taiji [19], LIGO [20] and pulsar timing arrays (PTAs) [21–23].

Big bang nucleosynthesis is highly sensitive to the Universe's expansion rate, as it determines the abundances of primordial elements like deuterium [24], helium [25], and lithium [26]. First-order PTs can modify the expansion history by altering the energy density and the equation of state. If such a transition occurs near or before BBN, it can impact element formation [27], allowing BBN observations to constrain PTs properties. CMB anisotropies are an important phenomenon in observation, since energy injection close to the recombination affects CMB anisotropies [28,29]. Additionally, CMB spectral distortions [30–35] that diverge the spectrum from the black body spectrum can

reveal changes in the energy distribution of photons. After considering the anisotropies from the first-order PT in the dark sector, Ref. [36] puts a restriction on the PT strength in the range of $10 - 10^6$ eV with the μ distortion in the CMB spectrum. Here, we focus on how the first-order PT modifies the energy evolution in the early Universe and affects the CMB spectrum.

In this paper, we investigate the effects of the low-scale first-order PTs whose temperatures are around sub-MeV and sub-keV, and place constraints on the PT parameters with BBN and CMB observational data, including the PT temperature, strength, and duration. We first briefly introduce the PTs' dynamics in Sec. II and discuss the BBN and CMB predictions from the low-scale slow first-order PTs in Secs. III and V. The conclusion is drawn in the last section.

II. PHASE TRANSITION

Considering first-order PTs, in the early Universe, the system was in the symmetric phase. As the Universe cooled, at the critical temperature T_c , the symmetric phase and the broken phase had the same free energy. With further cooling, true-vacuum bubbles of the broken phase nucleated within the symmetric phase, corresponding to the nucleation temperature T_n . Bubbles expand and percolate, eventually transforming the entire space into the new phase. The nucleation rate of the PT can be described as [37,38]

$$\Gamma(t) = \Gamma_0 e^{\beta t}, \quad (1)$$

where β determines the characteristic timescale of the PT, with the duration estimated as $\Delta t \sim \beta^{-1}$. In a radiation-dominated era, the prefactor can be approximated as

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$$\Gamma_0^{1/4} = \left(\frac{4\pi^3 g_\star}{45} \right)^{1/2} \left(\frac{T_p^2}{m_{\text{Pl}}} \right) e^{-\beta/8H_\star}, \quad (2)$$

where $T_p = T(t_p)$, and t_p is the percolation time, which is defined by $F(t_p) = 0.7$, where $F(t)$ implies the probability of the space remaining in the false vacuum state which will be explained below, and we refer t_p to PT time. $H_\star = H(T_p)$ represents the corresponding Hubble parameter. The Planck mass is given by $m_{\text{Pl}} \approx 1.22 \times 10^{19}$ GeV [39].

During the first-order PT, the evolution of the universe is governed by the Friedmann and Boltzmann equations [40,41]

$$H^2 = \frac{8\pi}{3m_{\text{Pl}}^2} (\rho_r + \rho_w + \rho_v), \quad (3)$$

$$\frac{d}{dt}(\rho_r + \rho_w) + 4H(\rho_r + \rho_w) = -\frac{d\rho_v}{dt}. \quad (4)$$

We assume the energy of the PT is injected into the plasma background. Here, the false vacuum energy density is expressed as $\rho_v = F(t)\Delta V$, with ΔV being the energy difference between the false and true vacuum. The quantities ρ_r and ρ_w represent the energy densities of radiation and bubble walls, respectively. The evolution of vacuum energy follows

$$-\frac{d\rho_v}{dt} = -\alpha\rho_r \cdot \frac{dF(t)}{dt}, \quad (5)$$

where the PT parameter $\alpha = \Delta V/\rho_r(t_p)$ characterizes how strongly the transition modifies the background energy.

Here, the definition of the parameter α does not include the entropy change during the PT. Compared to Ref. [27], we consider a noninstantaneous PT, where $\beta/H_\star$ characterizes the duration of the transition—the smaller the value of $\beta/H_\star$, the slower the PT. In the limit $\beta/H_\star \rightarrow \infty$, the transition becomes instantaneous. However, in Ref. [27], only an instantaneous PT is considered, with the PT energy assumed to be injected into the cosmic background instantaneously.

The radius of a bubble nucleated at time t' is given by

$$r(t, t') = \int_{t'}^t a^{-1}(t) dt, \quad (6)$$

where $a(t)$ is the scale factor and we take $v_{\text{wall}} = 1$. The function $F(t)$ follows [42]

$$F(t) = \exp \left[-\frac{4\pi}{3} \int_{t_i}^t dt' \Gamma(t') a^3(t') r^3(t, t') \right], \quad (7)$$

where t_i is the onset time of the PT. Initially, one has $F(t < t') = 1$ in the false vacuum. As the PT proceeds, $F(t)$ gradually decreases. The energy from the PT is

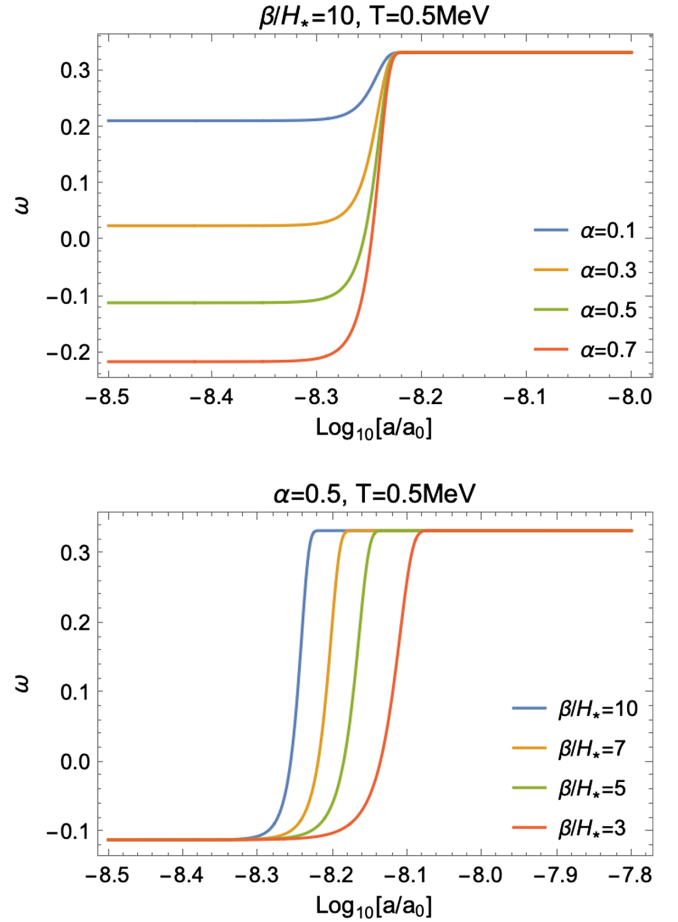


FIG. 1. The equation of state parameter ω versus the scale factor a for different PT strengths and inverse durations, a_0 is today's value.

transferred to the cosmic background, potentially causing reheating of photons or neutrinos.

The PT, which redistributes energy from ρ_v to ρ_r , can alter the Hubble expansion rate. We consider the total energy density $\rho = \rho_r + \rho_v$ and the total pressure $p = \rho_r/3 - \rho_v$, and compute the equation of state parameter as $\omega = p/\rho$. Figure 1 shows that stronger and slower PTs lead to larger deviations in ω , indicating a more significant modification of the expansion rate during the PT.

III. BIG BANG NUCLEOSYNTHESIS AND PTS

Neutrino decoupling is an important thermodynamic process in the early universe and is related to the parameter effective neutrino number, which is defined by $N_{\text{eff}} = 3 \times (11/4)^{4/3} (T_\nu/T_\gamma)^4$ at later times. To demonstrate the PT's effect on the value of $\Delta N_{\text{eff}} = N_{\text{eff}}^{\text{SM}} - N_{\text{eff}}$ when T_p is below 1 MeV, we present the case of photon reheating. In Fig. 2, we show the evolution of the photon energy density over time, highlighting the specific impact of a first-order PT on cosmic thermodynamics. During the PT, the photon energy density increases due to reheating, which alters the

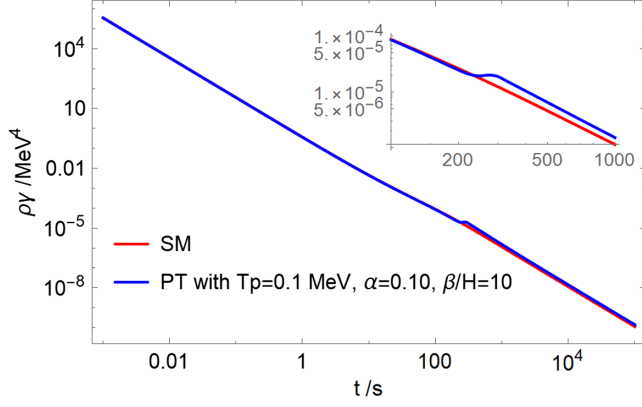


FIG. 2. The evolution of photon energy density over time. The red line corresponds to the standard model, while the blue line represents the case with a PT, where the PT parameter is $T_p = 0.1$ MeV, $\alpha = 0.1$, and $\beta/H_* = 10$.

neutrino-to-photon temperature ratio and consequently affects the value of ΔN_{eff} . By solving the time evolution of the neutrino and photon temperatures, we obtain the value of N_{eff} . More explicitly, the corresponding temperature evolution equations are [43],

$$\begin{aligned} \frac{dT_\gamma}{dt} &= - \left(4H\rho_\gamma + 3H(\rho_e + p_e) \right. \\ &\quad \left. + 3HT_\gamma \frac{dP_{\text{int}}}{dT_\gamma} + \frac{\delta\rho_{\nu_e}}{\delta t} + 2\frac{\delta\rho_{\nu_\mu}}{\delta t} + \frac{d\rho_{\text{vac}}}{dt} \right) / f(T_\gamma), \\ \frac{dT_\nu}{dt} &= - \left(12H\rho_\nu - \frac{\delta\rho_{\nu_e}}{\delta t} - 2\frac{\delta\rho_{\nu_\mu}}{\delta t} \right) / \left(3\frac{\partial\rho_\nu}{\partial T_\nu} \right). \end{aligned} \quad (8)$$

Here, $f(T_\gamma) = \frac{\partial\rho_\gamma}{\partial T_\gamma} + \frac{\partial\rho_e}{\partial T_\gamma} + T_\gamma \frac{d^2P_{\text{int}}}{dT_\gamma^2}$, and more details can be found in Ref [43]. We modified publicly available versions of `nuddec_BSM` [44,45] to take into account the PT dynamics and use them to compute the background thermodynamics.

Figure 3 mainly illustrates the impact of β/H_* on ΔN_{eff} , while also reflecting the effects of α and T_p . As long as $\beta/H_* > 1$, the PT can occur [46]. We chose a minimum value of $\beta/H_* = 10$ because, at that point, the impact on BBN is already significant. The upper two lines correspond to $\alpha = 0.01$, while the lower two correspond to $\alpha = 0.05$. The red and green lines represent $T_p = 0.1$ MeV, and the blue and purple lines represent $T_p = 0.5$ MeV. It can be observed that as α increases and β/H_* decreases, the impact of the PT becomes more significant, leading to a smaller N_{eff} . Compared to Ref. [27], the PT we consider has a stronger effect, allowing us to constrain α down to the order of 0.01.

We next briefly review the BBN process. Protons and neutrons can interconvert via weak reactions $n + \nu_e \leftrightarrow p + e^-$ and $n + e^+ \leftrightarrow p + \bar{\nu}_e$. At early times when the temperature is high, the weak reaction rate Γ_{np} exceeds the

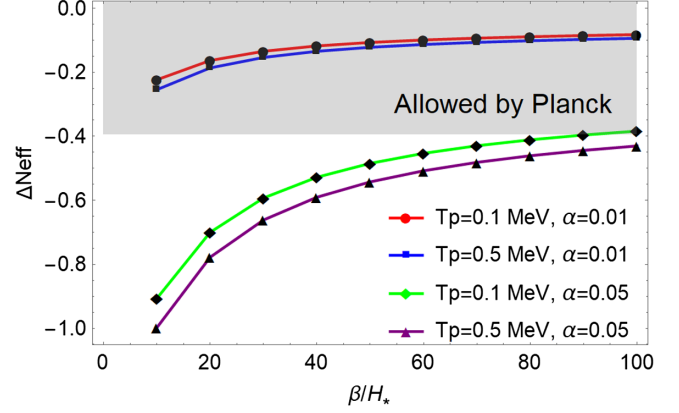


FIG. 3. The evolution of $\Delta N_{\text{eff}} = N_{\text{eff}} - N_{\text{eff}}^{\text{SM}}$ (with $N_{\text{eff}}^{\text{SM}} = 3.044(1)$ [45,47–49]) over β/H_* . The upper two lines illustrate the variation of the value of ΔN_{eff} with β/H_* under the conditions $\alpha = 0.01$. And the lower two lines correspond to $\alpha = 0.05$. The gray region shows the Planck 95% confidence level (CL) constraints $N_{\text{eff}} = 2.99^{+0.33}_{-0.34}$ [39].

Hubble expansion rate H , so protons and neutrons were maintained in thermal equilibrium. As the temperature dropped, the neutron fraction, defined as $X_n \equiv n_n/n_b$, gradually decreased, where n_x means the number density of x . The equilibrium neutron abundance is approximately given by $X_n^{\text{eq}} \approx e^{-Q/T}/(1 + e^{-Q/T})$, where $Q \equiv m_n - m_p = 1.293$ MeV. After weak interactions freeze out at temperature T_{FO} and time t_{FO} , X_n continues to decline slowly due to residual weak processes, and subsequently, free neutron decay $n \leftrightarrow p + e^- + \bar{\nu}_e$ becomes the dominant mechanism. During this epoch, the neutron fraction evolves as $X_n(t > t_{\text{FO}}) \approx X_n(t_{\text{FO}})e^{-(t-t_{\text{FO}})/\tau_n}$, where τ_n is the neutron lifetime. At this stage, the universe mainly consists of protons and neutrons. Deuterium (D) forms first, and only when deuterium exists in large amounts can helium be abundantly produced. The relative reactions are $n + p^+ \leftrightarrow D + \gamma$, $D + p^+ \leftrightarrow {}^3\text{He} + \gamma$, and $D + {}^3\text{He} \leftrightarrow {}^4\text{He} + p^+$. When the temperature reaches $T_{\text{nuc}} \approx 0.078$ MeV at time t_{nuc} , deuterium production becomes efficient, triggering the onset of rapid helium synthesis.

The final helium abundance $Y_P = \rho({}^4\text{He})/\rho_b$ is mainly determined by $X_n(t_{\text{nuc}})$, approximated as $Y_P = \rho({}^4\text{He})/\rho_b \approx 2X_n(t_{\text{nuc}})$. PTs occurring around the MeV scale can affect BBN [43], and those taking place below the MeV scale can also leave an impact. Considering the impact of PTs, variations in $X_n(t_{\text{FO}})$ and in the nucleosynthesis time t_{nuc} modify $X_n(t_{\text{nuc}})$, thereby directly affecting the final helium abundance. The relative importance of these two effects depends on the specific PT parameters. In scenarios with photon reheating from first-order PTs, the total energy density and the Hubble expansion rate are reduced at a given photon temperature $T = T_\gamma$. From Eq. (8), it can be seen that this causes the temperature to decrease more slowly with time, thereby delaying the onset of nucleosynthesis in

comparison with the SM, meaning that t_{nuc} is increased. This delay increases the fraction of neutrons that decay before nucleosynthesis, leading to a smaller helium abundance Y_p .

Conversely, the neutrino energy density is suppressed compared to the case without a PT, which lowers the weak interaction rate Γ_{np} at a given temperature. This can cause weak interactions to freeze out earlier, resulting in a higher neutron fraction $X_n(T_{\text{FO}})$. Although the reduced neutrino density also decreases the Hubble rate—partially mitigating the impact on Γ_{np} [50,51], the temperature dependence of Γ_{np} is stronger, so the net effect is a reduction in the reaction rate. Considering these two points together, the freeze-out temperature T_{FO} would be higher than in the scenario without a PT—that is, freeze-out occurs earlier, which likely yields a larger final Y_p .

Deuterium is abundantly produced first, and then further reacts to form helium. Its final abundance is primarily influenced by the timing of its destruction. Photon (neutrino) reheating from first-order PTs modifies the time-temperature relation, extending (shortening) the deuterium destruction period, resulting in a lower (higher) final $D/H|_p = n(D)/n(^1H)$. Considering the effects of PTs on the thermodynamic background, we modified the code PRIMAT [25] to calculate the corresponding nuclear reaction processes, thereby obtaining the abundances of helium and deuterium. And for a given isotope i , the evolution of

number density depends on the reaction rates in which it is involved. It is of the form

$$\frac{dn_i}{dt} + 3Hn_i = J_i, \quad (9)$$

where J_i is the net rate of evolution of the number density due to all nuclear reactions. There are more details in Ref. [25].

Figure 4 shows the effects of α and T_p on helium and deuterium abundances, where the Y_p and $D/H|_p$ are calculated with $\Omega_b h^2 = 0.021875$ and $\tau_n = 879.5$ s [52]. It can be observed that as α increases, no matter $T_p = 0.1$ MeV or 0.5 MeV both Y_p and $D/H|_p$ decrease for $\beta/H_* = 100$, while for $\beta/H_* = 10$, $T_p = 0.1$ MeV, helium abundance Y_p increases. Specifically, for $\beta/H_* = 100$, the influence of t_{nuc} dominates over that of T_{FO} , whereas for $T_p = 0.1$ MeV, $\beta/H_* = 10$, the opposite is true.

IV. CMB WITH FIRST-ORDER PT

In this section, we consider the effects of the PT on background energy transformation and, therefore, the CMB anisotropy, and also the CMB distortion driven by the additional energy injection from the PT.

A. Background energy distribution

The helium abundance Y_p plays a crucial role in determining the optical depth [53]

$$\kappa(\eta_L) \equiv \int_{t_L}^{t_0} \frac{dt}{\tau(t)} = \int_{\eta_L}^{\eta_0} \sigma_T n_t X_e a(\eta) d\eta. \quad (10)$$

This leads to greater suppression of CMB anisotropies, particularly at small scales, where $n_t = \frac{3H_0^2 \Omega_b}{8\pi G m_H} (1 - Y_p)$ is the present-day total electron number density including free and bound electrons, σ_T is Thomson scattering cross section, η is conformal time, $X_e = n_e/n_t$ is the free electron fraction, where $n_t = n_e + n_H$, n_e is the number density of free electron, and n_H is the number density of neutral hydrogen atom. Visibility function is given by

$$g = \frac{d\kappa(\eta_L)}{d\eta_L} \cdot e^{-\kappa(\eta_L)} = \sigma_T n_t X_e \cdot a(\eta_L) \cdot e^{-\kappa(\eta_L)}, \quad (11)$$

where a is the scale factor, η_L is the conformal time of last scattering. The visibility function reaches its maximum at the recombination epoch.

We utilize the publicly available cosmological code CLASS [54,55] to compute CMB anisotropy power spectra, the optical depth κ enters into the source function for CMB anisotropy [56]

$$S_{T0} = g \left(\frac{1}{4} \delta_\gamma + \psi \right) + e^{-\kappa} (\phi' + \psi'), \quad (12)$$

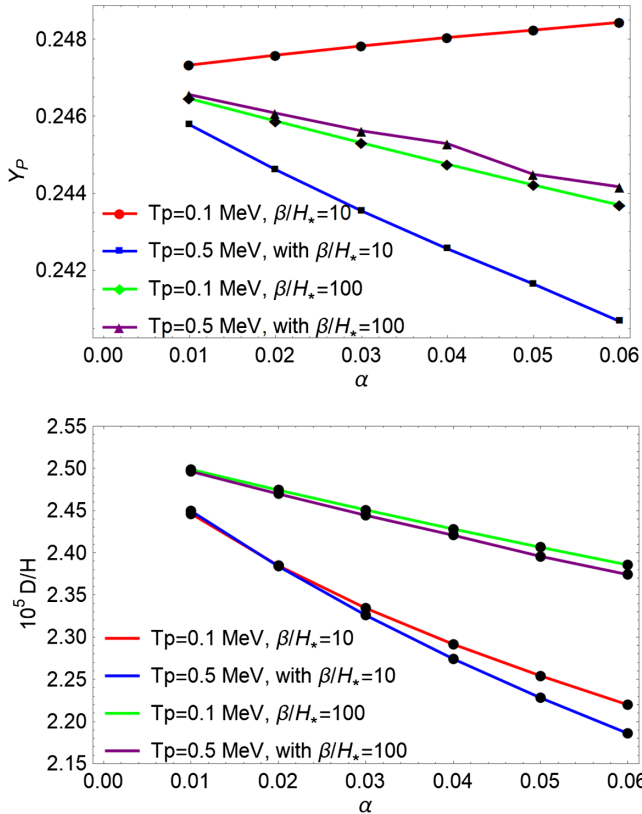


FIG. 4. The helium and deuterium abundances versus the PT strength α for different T_p and β/H_* .

where $\frac{1}{4}\delta_\gamma$ is intrinsic temperature fluctuation, ψ is gravitational redshift term, ϕ arises from gauge perturbations. The source function S_{T_0} is included in the photon transfer function

$$\Delta_\ell^{T_j}(q) = \int_{\tau_{\text{ini}}}^{\tau_0} d\tau S_{T_j}(k, \tau) \phi_\ell^j(\nu, \chi), \quad (13)$$

where ϕ_ℓ^j are radial functions linearly consisting of Bessel functions and their derivatives, this can be used to calculate the correlation function

$$C_\ell^{TT} = 4\pi \int \frac{dk}{k} \Delta_\ell^T(q, \tau_0) \Delta_\ell^T(q, \tau_0) \mathcal{P}(k), \quad (14)$$

where $\mathcal{P}(k)$ is the primordial spectrum, $q = \sqrt{k^2 + K}$ is the generalized wave number, with $K = -H_0^2(1 - \Omega_{\text{tot}})$.

We incorporated the PT mechanism into the background sector of the CLASS code, enabling energy transfer from vacuum energy to radiation. This modification is automatically linked to the BBN module, altering ΔN_{eff} . Here, we choose two benchmarks with PT temperatures of 0.5 MeV and 0.1 MeV to investigate how the PT impacts the CMB.

Figure 5 displays the angular power spectrum of correlation functions for temperature T anisotropies, with each

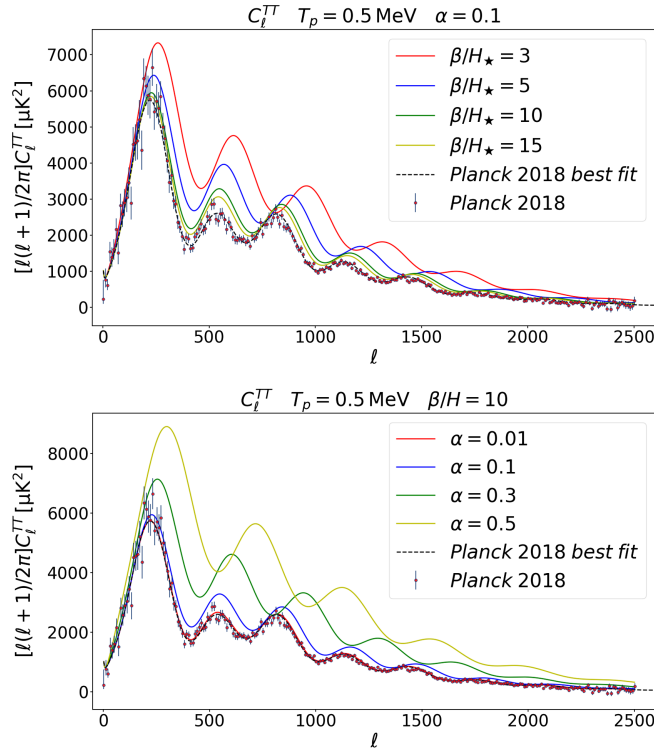


FIG. 5. The effects of PT parameters on the CMB TT spectrum, with the PT temperature being 0.5 MeV. The colored lines represent TT spectra with PT based on Planck 2018 best-fit parameters. The vertical lines with red dots represent the error bars from the Planck 2018 data [39].

curve representing different configurations of the PT parameters. Variations in the parameter α strongly affect the angular power spectrum, particularly around the peak, where large changes are observed. The Planck 2018 data is more restrictive on α compared to $\beta/H_\star$, with values typically around 0.1. The data also allow us to set a strong limit on the duration, requiring $\beta/H_\star \gtrsim 10$. The differences in spectra for the PT at 0.5 MeV and 0.1 MeV show an obvious pattern, as shown in Fig. 6. The difference grows as the $\beta/H_\star$ decreases and the α increases. This mainly arises from the influence of Y_p determined during the BBN epoch.

In Fig. 7, we further restrict the parameter space of PTs with the Planck 2018 data. We take 95% CL limits for $\chi^2 = \sum_l (C_l^{\text{PT}} - C_l^{\text{obs}})^2 / \sigma_l^{\text{obs}2}$, where σ_l^{obs} is $\pm 1\sigma$ diagonal uncertainty of Planck 2018. The plot highlights the constraints on the $\beta/H_\star$ and α , which weakly depend on the PT temperature, and PT parameter spaces with $\alpha \gtrsim \mathcal{O}(10^{-2})$ have been excluded for the slow PT at sub-MeV scales. Compared with the instantaneous PT scenario, which constrains $\alpha < 0.064$ for $T_p < 1$ MeV using the dataset $\{\Omega_b h^2, N_{\text{eff}}, Y_p\}$ from the baseline TTTEEE + low E analysis [27], our constraints on α are significantly tighter. Since we consider a durable PT scenario, we derive constraints on α and additional parameter $\beta/H_\star$ directly from the Planck 2018 observed data.

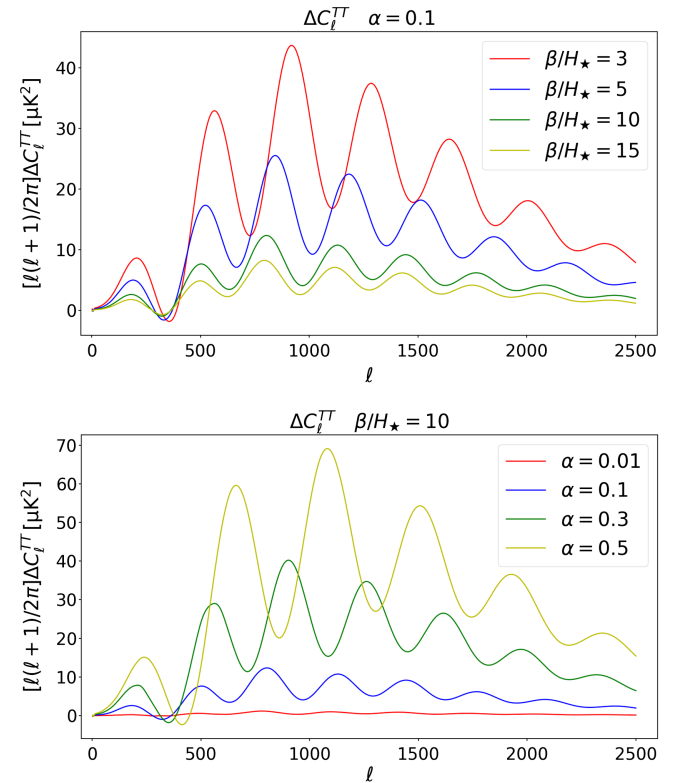


FIG. 6. The difference between the CMB anisotropy spectra for PTs at 0.5 MeV and 0.1 MeV, where $\Delta C_\ell^{TT} = C_\ell^{TT}(T_p = 0.5 \text{ MeV}) - C_\ell^{TT}(T_p = 0.1 \text{ MeV})$.

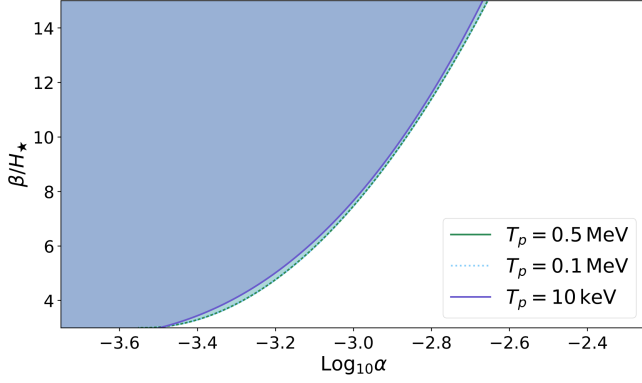


FIG. 7. The shaded regions represent the allowed parameter spaces for different PT temperatures by Planck 2018 data.

Our analysis employs a PT mechanism that differs from that of new early dark energy (NEDE) [29]. While NEDE models the PT as a source of early dark energy, we focus on a different scenario where vacuum energy is converted into radiation after the PT. In NEDE, the PT occurs very rapidly, characterized by $\beta/H_\star > 100$, with a corresponding redshift of $z_\star \sim 5000$. In contrast, our PT is significantly slower, with $\beta/H_\star \sim 10$, and occurs at a much higher redshift of approximately $z \sim 10^9$, close to the epoch of BBN. In our approach, CMB observational data constrain the parameters α and $\beta/H_\star$. These constraints ensure that the PT remains weak and occurs over a very short duration, effectively approximating an instantaneous transition. Our constraints suggest that a prolonged or strong transition would introduce deviations in the CMB anisotropy spectra, which are inconsistent with current observational data.

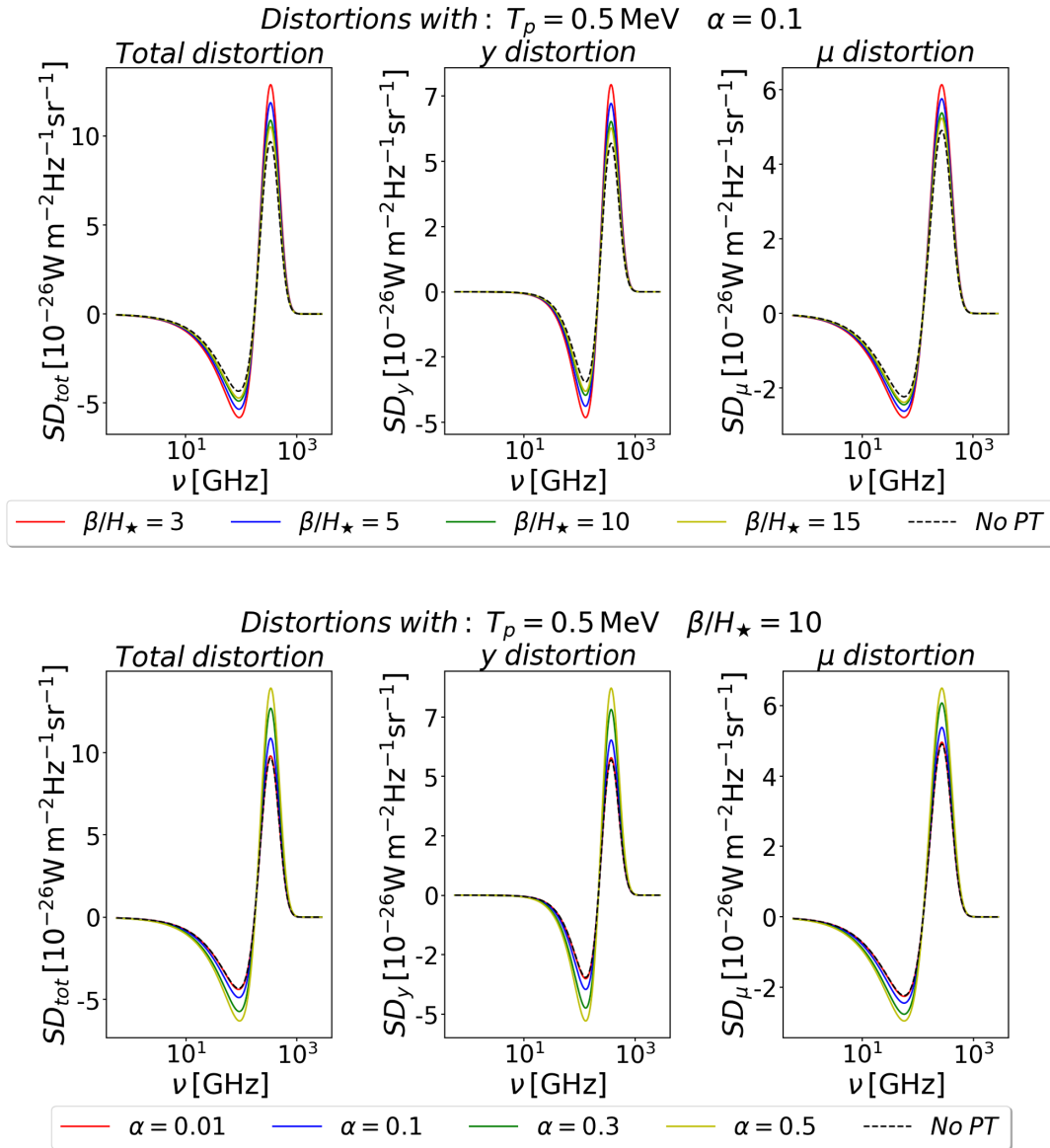


FIG. 8. The distortions of the noninjection scenario with varied PT parameters.

B. PT with CMB distortion

Next, we will discuss the effect of PTs on CMB spectral distortions. There are essentially two types of sources for spectral distortions: noninjection and injection.

Figure 8 illustrates the effect of noninjection energy redistribution on the noninjection spectral distortions. Each panel corresponds to a different set of PT parameters, showing how the redistribution of energy from ρ_v to ρ_γ modifies the resulting distortion. As can be seen, a larger α leads to a stronger spectral distortion, while a smaller β/H_* , corresponding to a longer-lasting PT, also enhances the distortion. We emphasize that this figure concerns the noninjection scenario, in contrast to the later discussion of sub-keV energy injection. Since the background energy redistribution in the noninjection scenario affects the entire evolution of the Universe, it can, in principle, occur at any PT temperature.

In contrast, energy injection acts as an additional source that does not modify the background evolution, and therefore it is restricted by the requirement on the PT temperature. At redshifts $z > 2 \times 10^6$, the evolution of photon number and energy density is highly efficient, maintaining the CMB spectrum as a blackbody. However, at lower redshifts $z < 2 \times 10^6$, the processes related to changes in photon number become less efficient, leading to the onset of spectral distortions. This period is characterized by efficient Compton scattering, which generates a chemical potential and is referred to as the μ -era. Following the μ -era, the y -era begins when Compton scattering also becomes inefficient, resulting in y -type spectral distortions. This era spans from $z_{\mu y} = 5 \times 10^4$ to the present. The μ -era begins to correspond to a cosmic temperature of approximately 500 eV. Spectral distortions are generated when energy injection occurs within this temperature range. Therefore, in this analysis, we focus on PTs below the keV scale, which is close to the μ -era.

Combining the heating rate $\dot{Q} = \frac{\partial \rho}{\partial t} + 4H\rho$ [31] with Eq. (4), we find the energy injection from PT to be

$$\dot{Q}_{\text{pt}} = -\frac{d\rho_v}{dt} = -\alpha\rho_r \cdot \frac{dF(t)}{dt}. \quad (15)$$

The μ distortion from the first-order PT is [31,57]

$$\mu \simeq 1.4 \int_{z_{\mu y}}^{\infty} dz \frac{dQ_{\text{pt}}}{dz} \frac{e^{-(z/z_{\mu})^5}}{\rho_r}, \quad (16)$$

where $z_{\mu} \simeq 1.98 \times 10^6$, and $\exp[-(z/z_{\mu})^5]$ is the visibility function that removes contributions to μ distortions at $z > z_{\mu}$. The transformation between redshift z and time t is given by $dt = -dz/((1+z)H(z))$.

The COBE/FIRAS experiment [58] sets an upper limit of $|\mu| < 9 \times 10^{-5}$ based on observational data. On the other

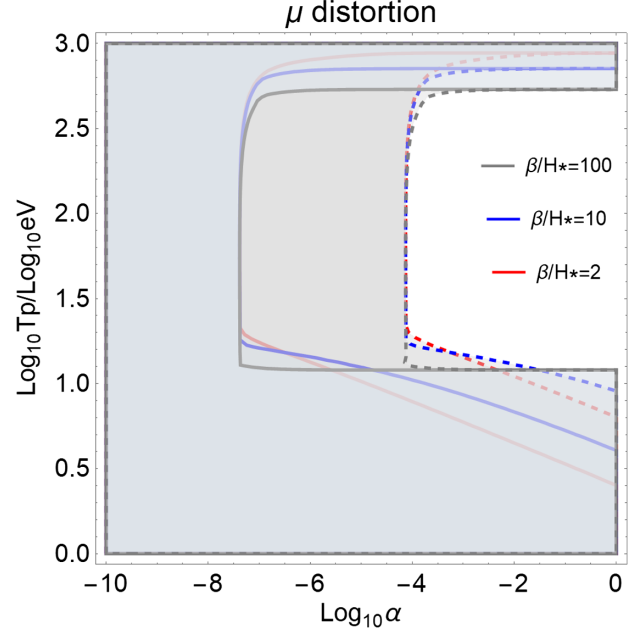


FIG. 9. The μ distortion versus different PT parameters. The shaded area denotes the allowed parameter space, while the dashed and solid lines correspond to $|\mu| = 9 \times 10^{-5}$ [58] and $|\mu| = 5 \times 10^{-8}$ [59], respectively.

hand, the proposed PIXIE mission [59], though not realized, was designed with a sensitivity goal of $|\mu| < 5 \times 10^{-8}$. In Fig. 9, we illustrate the impact of the PT parameters—the PT strength α [ranging from $\mathcal{O}(10^{-10})$ to $\mathcal{O}(1)$] and the PT temperature T_p (ranging from eV to keV)—on the μ distortion. The PT with the duration coincides with the time range in which μ -type distortions are generated, leading to similar levels of energy injection. The magnitude of the μ distortion corresponds to the PT strength for T_p between roughly $\mathcal{O}(10)$ eV and $\mathcal{O}(10^2-10^3)$ eV. And, for low PT temperature scenarios with $T_p \lesssim \mathcal{O}(10)$ eV, more parameter spaces are allowed for a faster PT with larger β/H_* and weaker PT with smaller α .

Compared with the first-order PTs in a dark sector, where the PT energy is transferred into sound waves in the dark fluid and SM plasma are considered [36], our study focuses on the direct heating of the photon-baryon fluid below the keV scale. The allowed range of the dimensionless parameter α , constrained in our analysis, shows a strong dependence on T_p , spanning from $\mathcal{O}(10^{-10})$ up to $\mathcal{O}(1)$.

V. GRAVITATIONAL WAVE

During the first-order PT, expanding bubbles of the new phase collide and transfer energy to the surrounding plasma, creating motion in the form of sound waves. The GW spectrum of the sound waves is characterized by the strength of the PT α , the inverse duration β/H , the

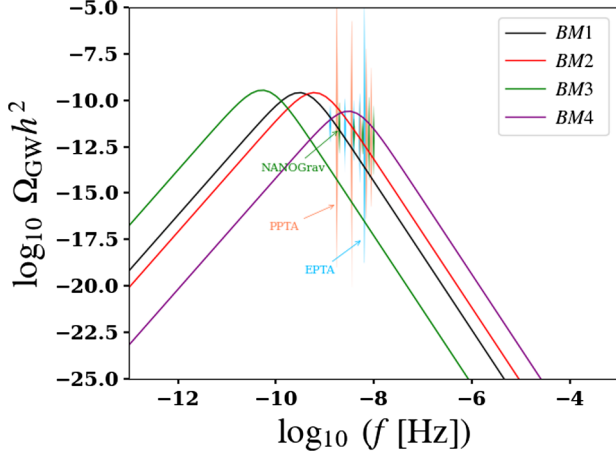


FIG. 10. We show the results from the four benchmark points listed in Table I. Solid lines show GW spectra from sound waves with different parameter choices, where $v_w = 1$, $\tau_{sw}H_\star = 1$, $\kappa_{sw} = 0.2$. We also present the sensitivities of different GW detectors, i.e., NANOGrav [21], PPTA [22], and EPTA [23].

PT temperature, and the bubble wall velocity. We use GW spectrum given in Refs. [21,60–63]

$$\begin{aligned} \Omega_{\text{GW}}^{\text{sw}} h^2(f) = & 7.7 \times 10^{-5} g_*^{-\frac{1}{3}} \left(\frac{\kappa_{\text{sw}} \alpha_*}{1 + \alpha_*} \right)^2 \left(\frac{H_*}{\beta} \right) 0.513 v_w \\ & \times (1 - (1 + 2\tau_{\text{sw}} H_*)^{-\frac{1}{2}}) (f/f_p^{\text{sw}})^3 \\ & \times [7/(4 + 3(f/f_p^{\text{sw}})^2)]^{\frac{7}{2}}, \end{aligned} \quad (17)$$

where $f_p^{\text{sw}} = 1.13 \times 10^{-10} \text{ Hz} \frac{0.536}{v_w} \left(\frac{\beta}{H_*} \right) \left(\frac{T_*}{\text{MeV}} \right) \left(\frac{g_*}{10} \right)^{1/6}$, v_w is bubble wall velocity, κ_{sw} is fraction of vacuum energy transform into sound wave, $\tau_{\text{sw}} H_\star$ is about lifetime of sound wave.

As illustrated in the Fig. 10, the resulting GW signals from the PTs with $0.1 \lesssim \alpha$, $10 \lesssim \beta/H_\star \lesssim 100$ at sub-MeV scales $0.1 \text{ MeV} \lesssim T_\star \lesssim 1 \text{ MeV}$ fall within the sensitivity range of PTAs' observations. Therefore, the sub-MeV scale slow first-order PT can be complementarily constrained by GW detectors and CMB anisotropy spectra. Ref. [64] shows that the slow first-order PT in the dark sector with subsequently decays before the time of neutrino decoupling is a compelling possibility to account for the signatures in the NanoGrav 12.5-yr dataset. The chance to address dark matter and PTAs' observables with a sub-GeV scale first-order PT in the dark sector has been explored in Ref. [65].

TABLE I. Benchmark points in Fig 10.

	α	$\beta/H_\star$	$T_\star(\text{MeV})$
BM_1	0.1	10	0.5
BM_2	0.1	10	1
BM_3	0.1	10	0.1
BM_4	0.1	100	0.5

VI. CONCLUSION AND DISCUSSION

In this study, we explore the possibility of a low-scale first-order PTs occurring around or after the epoch of BBN. We find that ΔN_{eff} can constrain the PT strength α down to the order of 0.01, and nearly all cases with $\alpha > 0.05$ are excluded by Planck observables data for slow PTs with $\beta/H_\star < 100$. The modification of the background energy evolution induced by the PTs, combined with the altered BBN yields, imposes effects on the CMB anisotropies. The differences in anisotropy spectra show a regular pattern, at the level of $\mathcal{O}(10)$. For sub-MeV scale PTs, we find that the CMB anisotropy imposes tight constraints, requiring the PT strength to satisfy $\alpha \lesssim 0.01$ and its inverse duration to meet $\beta/H_\star \gtrsim 5$.

For PTs below the keV scale, energy injection can induce μ distortions. From this effect, we obtain constraints on T_p and α in the range from $\mathcal{O}(10^{-10})$ to $\mathcal{O}(1)$ for slow PTs. Additionally, PTAs constrain strong PTs with $\alpha \gtrsim \mathcal{O}(10^{-1})$ in the regime of slow transitions, specifically for $\beta/H_\star \lesssim 10$.

In addition to the energy evolution effects studied here, first-order PTs can also leave imprints through the generation of perturbations. Both Refs. [40] and [66] investigate super-horizon curvature perturbations induced by late-time first-order cosmological PTs. Ref. [40] derives constraints for PTs with temperatures in the range $T_\star \in [10^{-6}, 10^3] \text{ GeV}$, $\alpha \in [10^{-4}, 1]$ for slow PTs with $\beta/H_\star \in [1, 100]$. They show that for the PTs with $\beta/H_\star \leq 10$, the parameter space $\alpha \geq 2 \times 10^{-3}$ has been excluded model independently for $T_\star \in [0.3, 100] \text{ MeV}$. And Ref. [66] considers the effect of the finite bubbles statistics, obtains the stringent constraints on the parameters of such late-time PTs, which even surpass the ΔN_{eff} limits when $0.1 \text{ eV} \leq T_{\text{PT}} \leq 1 \text{ keV}$. These results suggest that the induced perturbations may leave distinct imprints in the CMB and matter power spectra, posing challenges to some models proposed to resolve the Hubble tension. Reference [67] shows that tensor perturbations produced by a late-time around the eV and keV, strongly first-order cosmological PT can also generate a B-mode signal, whose strength can rival testable inflationary predictions across a range of observable scales. The PTs exhibit maximal power on small scales $k \geq 0.01 \text{ Mpc}^{-1}$, and become suppressed on large scales.

ACKNOWLEDGMENTS

This work is supported by the National Key Research and Development Program of China under Grant No. 2021YFC2203004, and by the National Natural Science Foundation of China (NSFC) under Grants No. 12322505, No. 12347101. We also acknowledge Chongqing Talents: Exceptional Young Talents Project No. cstc2024ycjh-bgzxm0020 and Chongqing Natural Science Foundation under Grant No. CSTB2024NSCQ-JQX0022.

DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

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