New window into gravitationally produced scalar dark matter

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Conventional scenarios of purely gravitationally produced dark matter with masses below the Hubble parameter at the end of inflation are in tension with cosmic microwave background (CMB) constraints on the isocurvature power spectrum. We explore a more general scenario with a nonminimal coupling between the scalar dark matter field and gravity, which allows for significantly lighter scalar dark matter masses compared to minimal coupling predictions. By imposing relic abundance, isocurvature, Lyman-α, and big bang nucleosynthesis (BBN) constraints, we show the viable parameter space for these models. Our findings demonstrate that the presence of a nonminimal coupling expands the parameter space, yielding a dark matter mass lower bound of 2×10^{-4} eV.

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I. INTRODUCTION

The nature and origin of dark matter (DM) remain one of the greatest unsolved mysteries in fundamental physics. Furthermore, the lack of detection from indirect DM and terrestrial experiments, coupled with the stringent limitations imposed by direct DM detection searches such as XENON1T [1], LUX [2], PandaX [3], and LZ [4], challenges the conventional weakly interacting massive particle (WIMP) paradigm without providing new insights into the composition of the universe's invisible component. This discrepancy necessitates exploring alternative DM models [5–7].

One of the most well-motivated and notably minimalistic models involves the gravitational production of hidden sector particles during the transition from the inflationary quasi-de Sitter phase to a matter-dominated (MD) or radiation-dominated (RD) universe [8–12]. During reheating, inflaton oscillations lead to additional particle production [13–15]. However, for nearly 20 years, it has been known that the CMB bounds on the amplitude of the isocurvature power spectrum imply that the purely gravitationally produced DM must be superheavy, i.e., close to the Hubble scale at the end of inflation [16–22]. In this work, we demonstrate that when the spectator DM field couples nonminimally to gravity [23–28], such models avoid the isocurvature constraints [16,23], opening up the parameter space even for light DM with masses significantly below the Hubble scale at the end of inflation.

II. FRAMEWORK

We consider the homogeneous and isotropic Friedmann-Robertson-Walker (FRW) metric $\mathrm{d}s^2 = a(\eta)^2(\mathrm{d}\eta^2 - \delta_{ij}\mathrm{d}x^i\mathrm{d}x^j)$, where $a(\eta)$ represents the scale factor and $\mathrm{d}\eta = \mathrm{d}t/a$ is the conformal time. The general action of our theory is given by

$$\mathcal{S} = \int \mathrm{d}^4x\sqrt{-g}\bigg[-\frac{1}{2}(M_P^2 - \xi\chi^2)R + \frac{1}{2}(\partial_\mu\phi)^2 - V(\phi) + \frac{1}{2}(\partial_\mu\chi)^2 - \frac{1}{2}m_\chi^2\chi^2\bigg]. \quad (1)$$

Here $g = \det g_{\mu\nu}$ represents the determinant of the metric, $M_P = 1/\sqrt{8\pi G_N} \simeq 2.435 \times 10^{18}$ GeV denotes the reduced Planck mass, R is the Ricci scalar, ξ is the nonminimal coupling of the DM field to gravity, with $\xi = 0$ and $\xi = 1/6$ corresponding to minimal and conformal couplings, respectively. ϕ is the inflaton field, where $V(\phi)$ is the corresponding potential, and χ is the spectator scalar DM field whose bare mass is denoted by m_χ.

Introducing the canonically normalized field $X \equiv a\chi$, and varying the action (1) with respect to X, we obtain the equation of motion

$$(\partial_\eta^2 - \nabla^2 + a^2 m_{\mathrm{eff}}^2)X = 0, \quad m_{\mathrm{eff}}^2 = m_\chi^2 + \frac{1}{6}(1-6\xi)R. \quad (2)$$

During inflation, one can approximate the Ricci scalar as equal to its de Sitter value, $R \simeq -12H^2$, with H the Hubble parameter, and the effective mass with minimal coupling ($\xi = 0$) becomes $m_{\mathrm{eff}}^2 = m_\chi^2 - 2H^2$, whereas with conformal coupling ($\xi = 1/6$), it becomes $m_{\mathrm{eff}}^2 = m_\chi^2$. This implies that light scalars $m_\chi \ll H$ minimally coupled to gravity would experience a tachyonic phase during inflation with $m_{\mathrm{eff}}^2 < 0$. Problematically, during inflation, the tachyonic growth of light DM modes can efficiently generate isocurvature perturbations at the second order and lead to an unsuppressed isocurvature spectrum for long-wavelength (IR) modes [16,17,21,29], in disagreement with the current isocurvature power spectrum constraints from *Planck* observations [30]. Nevertheless, when the conformal coupling is sufficiently large, one can approximate the effective mass as $m_{\mathrm{eff}}^2 \simeq 12\xi H^2$, which implies that the DM effective mass is very large during inflation but becomes much smaller at the end of inflation when the Hubble parameter drops to lower values, successfully avoiding the current isocurvature bound.

Since the FRW metric is spatially homogeneous, one can perform a Fourier expansion of the DM field X:

$$X(\eta, \boldsymbol{x}) = \int \frac{\mathrm{d}^3\boldsymbol{k}}{(2\pi)^{3/2}} e^{-i\boldsymbol{k}\cdot\boldsymbol{x}} \left[X_k(\eta)\hat{a}_{\boldsymbol{k}} + X_k^*(\eta)\hat{a}_{-\boldsymbol{k}}^\dagger\right], \quad (3)$$

where $\boldsymbol{k}$ is the comoving momentum, with $|\boldsymbol{k}| = k$, and $\hat{a}_{\boldsymbol{k}}$ and $\hat{a}_{\boldsymbol{k}}^\dagger$ are the annihilation and creation operators, respectively, which obey the canonical commutation relations $[\hat{a}_{\boldsymbol{k}}, \hat{a}_{\boldsymbol{k}'}^\dagger] = \delta(\boldsymbol{k} - \boldsymbol{k}')$ and $[\hat{a}_{\boldsymbol{k}}, \hat{a}_{\boldsymbol{k}'}] = [\hat{a}_{\boldsymbol{k}}^\dagger, \hat{a}_{\boldsymbol{k}'}^\dagger] = 0$. The canonical commutation relations between the field, X_k, and its momentum conjugate, X_k', are satisfied if the Wronskian condition $X_k X_k^{*\prime} - X_k^* X_k' = i$ is imposed. If we substitute the Fourier decomposed field (3) into the equation of motion (2), we find that the equation of motion together with the angular frequency are given by

$$X_k'' + \omega_k^2 X_k = 0, \quad \text{with} \quad \omega_k^2 = k^2 + a^2 m_{\mathrm{eff}}^2. \quad (4)$$

III. GRAVITATIONAL PRODUCTION OF DM

To compute the gravitational production of DM during inflation and reheating stages, one must specify the initial conditions and solve the mode equations (4). In the early-time asymptotic limit $\eta \to -\infty$, modes deep inside the horizon $aH \ll k$ satisfy $k^2/a^2 \gg R$ and the mode frequency can be approximated as $\omega_k \simeq k$. This motivates the choice of the Bunch-Davies vacuum initial condition, given by $\lim_{\eta\to-\infty} X_k(\eta) = \frac{1}{\sqrt{2k}} e^{-ik\eta}$. The comoving number density of gravitationally produced scalar DM, n_χ, can be computed using the following expression [21,31,32]:

$$n_\chi \left(\frac{a}{a_{\mathrm{end}}}\right)^3 = \int_{k_0}^{\infty} \frac{\mathrm{d}k}{k} \mathcal{N}_k \quad \text{with} \quad \mathcal{N}_k = \frac{k^3}{2\pi^2} f_\chi(k,t), \quad (5)$$

where a_{end} is the scale factor at the end of inflation, and $\mathcal{N}_k$ is the comoving number density spectrum expressed as a function of the DM phase space distribution, $f_\chi(k,t) = \frac{1}{2\omega_k}|\omega_k X_k - iX_k'|^2$. Here, we have introduced an IR cutoff, $k_0 = a_0 H_0$, where $a_0 = a(\eta_0)$ is the present-day scale factor and $H_0 = H(\eta_0)$ represents the present-day Hubble parameter. We note that k_0 represents the present comoving scale, assuming that inflation started when this mode was inside the horizon. Modes with lower wave numbers are outside of our cosmological horizon and, as a result, contribute to the homogeneous background [33]. We emphasize that our results are insensitive to the IR cutoff scale. The isocurvature constraints discussed later restrict the allowed parameter space to be only sensitive to the UV. A detailed discussion regarding the significance of the IR cutoff can be found in Ref. [29]. Note that both tachyonic instabilities and resonance effects are accounted for by solving the mode equation (4) in the presence of the background induced by the inflaton field during and after inflation.

It can be shown both analytically and numerically that the phase space distribution in the long-wavelength (IR) regime scales as $f_\chi \propto k^{-2\nu}$ ($\mathcal{N}_k \propto k^{3-2\nu}$), where $\nu = \sqrt{9/4 - 12\xi - m_\chi^2/H^2}$ for real ν. If we assume that the DM scalar is light, with $m_\chi \ll H$, we find that $f_\chi \propto k^{-3}$ ($\mathcal{N}_k = \text{const.}$, flat spectrum) for minimal coupling and $f_\chi \propto k^{-1}$ ($\mathcal{N}_k \propto k^2$) for conformal coupling [13,19,21,34]. The comoving number density (5) contains an IR divergence for $\xi = 0$, which is regulated by k_0. However, the integral is convergent for $\xi > 0$ and becomes IR insensitive approximately when $\xi \gtrsim 1/5$.

We illustrate the characteristic dependence of the comoving number density spectrum on ξ in Fig. 1. We note that $\mathcal{N}_k$ typically peaks at $k_{\mathrm{peak}} > k_{\mathrm{end}}$, where $k_{\mathrm{end}} = a_{\mathrm{end}} H_{\mathrm{end}}$ is the mode that reenters the horizon at the end of inflation. We evaluate the comoving number density at late time $a_{\mathrm{final}}/a_{\mathrm{end}} = 100$ when the distribution stops evolving and there is no additional particle production, ensuring that $n_\chi \left(\frac{a}{a_{\mathrm{end}}}\right)^3 = \text{const.}$

The mode k_{peak}, and the short-wavelength (UV) tail of the spectrum, correspond to modes that remain inside the horizon during inflation. For $k > k_{\mathrm{preak}}$, the spectrum scales as $f_\chi \propto k^{-9/2}$ ($\mathcal{N}_k \propto k^{-3/2}$), independently of the value of ξ. The amplitude of $\mathcal{N}_k$ at k_{peak} increases with increasing value of the coupling ξ.

For a large coupling ξ, the effective mass (2) increases, along with the parameter ν, and the spectrum becomes

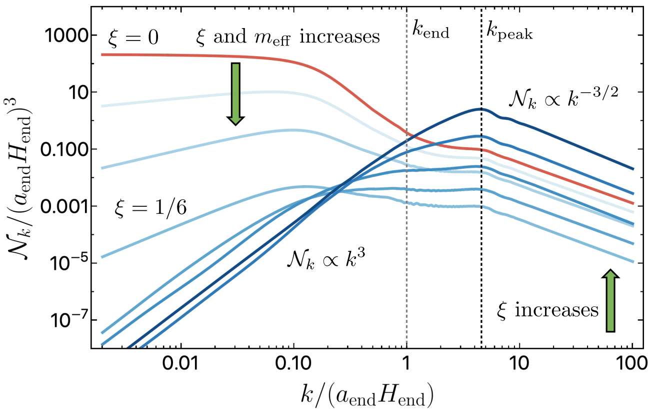

FIG. 1. A qualitative diagram illustrating the dependence of the comoving number density spectrum $\mathcal{N}_k$ on nonminimal coupling ξ as a function of rescaled horizon modes $k/(a_{\rm end}H_{\rm end})$.

suppressed in the IR. The same effect occurs in the case of direct DM-inflaton coupling [13,29]. When the coupling ξ is significant, $m_{\rm eff} \simeq 12\xi H^2$, and the transition from quasi-dS to a matter- or radiation-dominated universe causes the effective mass to change rapidly, leading to substantial particle production. Furthermore, when ξ becomes large (and ν is imaginary), we find that $f_\chi = \text{const.}$ ($\mathcal{N}_k \propto k^3$).

Assuming that reheating occurs significantly after the end of inflation, with $a_{\rm reh} \gg a_{\rm end}$, and that the total entropy $S = sa^3$ is conserved after the end of reheating, and using Eq. (5), we find that the DM relic density can be expressed as [13]

$$\Omega_{\rm DM,grav} \simeq \frac{m_\chi n_\chi}{\rho_c} = \frac{1}{6\pi q_0^3}\left(\frac{m_\chi H_0}{M_P^2}\right)\int_{q_0}^{\infty} {\rm d}q\, q^2 f_\chi(q). \quad (6)$$

Here, $q \equiv k/(a_{\rm end}H_{\rm end})$ is the rescaled dimensionless comoving momentum, with $q(k_{\rm end}) = 1$, $H_0 = 100h\ {\rm km\, s^{-1}\, Mpc^{-1}}$ is the present Hubble parameter, $\rho_c = 1.05\times 10^{-5}h^2\,{\rm GeV\, cm^{-3}}$ is the present critical energy density, with $h \sim 0.67$ [35], and the present comoving scale (IR cutoff) is given by $q_0 \simeq 7\times 10^{-30}(M_P^2/(H_{\rm end}T_{\rm reh}))^{1/3}$ [13,36–38]. We must ensure that the measured DM relic abundance satisfies the experimental value $\Omega_{\rm DM}h^2 = 0.1198$ [35], and it is obtained at some reheating temperature in the range $T_{\rm BBN} \le T_{\rm reh} \le T_{\rm max}$ with the BBN temperature of $T_{\rm BBN} \simeq 1$ MeV and $T_{\rm max} = (90H_{\rm end}^2 M_P^2/\pi^2 g_{\rm reh})^{1/4}$, which is the theoretical maximum temperature, corresponding to the instantaneous conversion of the total inflaton energy density to radiation at the end of inflation.

IV. RESONANT DM PRODUCTION

For large nonminimal coupling $\xi \gtrsim 10$, the mode equation (4) becomes a Mathieu equation, which features parametric instabilities, causing exponential quasi-stochastic excitations of modes $q \sim 1$–10, as discussed further in the Supplementary Material [39]. Full numerical solutions are needed to account for these effects, which we consider up to $\xi \simeq 70$. Beyond that, the DM energy density exceeds one percent of the inflaton energy density, requiring more advanced tools, such as lattice simulations, as explored in Ref. [40,41].

V. THERMAL PRODUCTION OF DM

One must account for the thermal production of DM, which occurs through the gravitational scattering of Standard Model (SM) particles from the thermal bath is given by

$$R_\chi = \frac{\pi^3(2560\xi(3\xi-1)+3997)}{41472000}\frac{T^8}{M_P^4}, \quad (7)$$

and the thermally-produced DM relic abundance is

$$\Omega_{\rm DM,thermal} \simeq 1.9\times 10^9 g_{\rm reh}^{-3/2}\left(\frac{M_P R_\chi(T_{\rm reh})}{T_{\rm reh}^5}\right)\left(\frac{m_\chi}{1\ {\rm GeV}}\right). \quad (8)$$

The detailed calculations are given in the Supplemental Material (SM).

VI. ISOCURVATURE CONSTRAINTS

When a light scalar DM field (spectator field) is excited during inflation, it inevitably leads to large isocurvature perturbations [16,17]. The rapid increase in DM energy density is primarily driven by the quadratic fluctuations that substantially contribute to the variance $\langle\chi^2\rangle$ [16,18,21,22]. Our analysis assumes no initial misalignment for the DM scalar field at the beginning of inflation, with $\langle\chi\rangle = 0$, (and $\langle\chi^2\rangle = 0$) [29,33]. Notably, when $\langle\chi\rangle = 0$, the DM inhomogeneities do not directly affect the curvature perturbation, and they can be treated as pure isocurvature fluctuations in the comoving gauge [16,17]. The second-order contribution to the isocurvature power spectrum is given by [16,21,42]

$$\mathcal{P}_\mathcal{S}(k) = \frac{k^3}{2\pi^2\rho_\chi^2}\int {\rm d}^3\boldsymbol{x}\,\langle\delta\rho_\chi(\boldsymbol{x})\delta\rho_\chi(0)\rangle e^{-i\boldsymbol{k}\cdot\boldsymbol{x}}, \quad (9)$$

where ρ_χ and $\delta\rho_\chi$ denote the DM energy density and its fluctuation, respectively. The current constraints on the isocurvature power spectrum provided by *Planck* are $\beta_{\rm iso} \equiv \mathcal{P}_\mathcal{S}(k_*)/(\mathcal{P}_\mathcal{R}(k_*)+\mathcal{P}_\mathcal{S}(k_*)) < 0.038$ at the 95% confidence limit (C.L.) for the pivot scale $k_* = 0.05\ {\rm Mpc}^{-1}$, where $\mathcal{P}_\mathcal{R}(k_*) = 2.1\times 10^{-9}$ is the curvature power spectrum [30]. This imposes an upper limit on the isocurvature power spectrum $\mathcal{P}_\mathcal{S}(k_*) \lesssim 8.3\times 10^{-11}$.

Our results rely on a numerical evaluation of Eq. (9). For a more intuitive understanding of the isocurvature constraint, one can assume a constant effective mass $m_{\rm eff}/H \lesssim 1$ with a constant Hubble parameter. In this case, the mode function on super-horizon scales can be approximated by

$$|\chi_k(k \ll aH)| \simeq \frac{H}{\sqrt{2a^3H^3}}\left(\frac{k}{aH}\right)^{-3/2+m_{\rm eff}^2/(3H^2)}. \qquad (10)$$

Such super-horizon modes are exponentially suppressed by the total duration of inflation in number of e-folds ΔN spent outside the horizon $|\chi_k(k \ll aH)| \sim e^{-\Delta N m_{\rm eff}^2/(3H^2)}$. Such long-wavelength contribution can be regarded as coherent oscillations of the DM field [11,43] with a displacement at the end of inflation is given by $\sqrt{\langle\chi^2\rangle}$ with $\langle\chi^2\rangle \sim H^4/m_{\rm eff}^2$. Approximating the field energy-density by $\rho_\chi \simeq m_{\rm eff}^2\langle\chi^2\rangle/2$, as the mode-function convolution in the numerator of Eq. (9) inherits the exponential ΔN suppression, the isocurvature power spectrum scales as $\mathcal{P}_\mathcal{S}(k_*) \sim (m_{\rm eff}^4/H^4)e^{-4\Delta N_* m_{\rm eff}^2/(3H^2)}$, which becomes more suppressed as ξ, and in turn $m_{\rm eff}$, increases. Details of our derivation can be found in the Supplementary Material sections.

For purely gravitational DM production, the isocurvature constraints from *Planck* require that $m_\chi^2 + 12\xi H_*^2 \gtrsim H_*^2/4$, where H_* is the Hubble scale at the horizon exit. For a minimal coupling ($\xi = 0$), this constraint becomes $m_\chi \gtrsim H_*/2$ [29]. Assuming a very light bare DM mass, with $m_\chi \ll 2\sqrt{3\xi}H_*$, we find that the isocurvature limits are always avoided for $\xi > 1/48$. This implies that even a very small value of a nonminimal coupling is sufficient to satisfy *Planck* isocurvature limits, and the conformal coupling case ($\xi = 1/6$) always satisfies the constraint.

VII. LYMAN-α FOREST CONSTRAINTS

In contrast to the conventional WIMP scenario, light DM particles produced in an out-of-equilibrium state may possess a considerable pressure component, resulting in the suppression of overdensities on galactic scales and subsequently implying a cutoff in the matter power spectrum $\mathcal{P}(k)$ for scales k larger than the free-streaming horizon wavenumber $k_{\rm H}(a)$. Lyman-α forest measurements constrain such cutoff scale. Limits are customarily expressed in terms of a lower bound on the mass of a generic warm DM (WDM) candidate decoupling from the thermal bath $m_{\rm WDM} > m_{\rm WDM}^{\text{Ly-}\alpha} \simeq (1.9-5.3)$ keV at 95% C.L. [44–50], which corresponds to a cutoff scale of the order of $k_{\rm H}(a=1) > 15h$ Mpc^{-1} [51]. For DM produced in an out-of-equilibrium state, one can translate the Lyman-α bound into a constraint on the DM mass by matching the corresponding equation of state parameters. This process leads to the following bound [13,51]

$$m_\chi^{\text{Ly-}\alpha} = m_{\rm WDM}^{\text{Ly-}\alpha}\left(\frac{T_\star}{T_{\rm WDM,0}}\right)\sqrt{\frac{\langle q^2\rangle}{\langle q^2\rangle_{\rm WDM}}}, \qquad (11)$$

where $T_{\rm WDM,0}$ is the WDM temperature saturating the DM abundance and this quantity only depends on $m_{\rm WDM}^{\text{Ly-}\alpha}$ [51]. $T_\star = H_{\rm end}(a_{\rm end}/a_0)$ is the characteristic energy scale of the produced DM, and the normalized second moment of the DM phase space distribution function $\langle q^2\rangle \equiv \int dq\, q^4 f_\chi(q)/\int dq\, q^2 f_\chi(q)$. For a WDM candidate, this quantity is approximately $\langle q^2\rangle_{\rm WDM} \simeq 12.93$. We argue that large values of nonminimal coupling ξ allow the scalar DM to be very light, and the values of ξ can be constrained by the structure formation limits, as detailed in our parameter space analysis below.

VIII. RESULTS AND DISCUSSION

To impose the model constraints and demonstrate the effect of a nonminimal coupling ξ, we consider the T-model inflationary potential [52]

$$V(\phi) = \lambda M_P^4\left[\sqrt{6}\tanh\left(\frac{\phi}{\sqrt{6}M_P}\right)\right]^2, \qquad (12)$$

where the potential can be normalized using the approximation $\lambda \simeq 3\pi^2 A_{S*}/N_*^2$ [53]. For a nominal choice of $N_* = 55$ e-folds, we obtain $\lambda \simeq 2\times10^{-11}$, the spectral tilt $n_s \simeq 0.963$, and the tensor-to-scalar ratio $r \simeq 0.004$, which is highly favored by current CMB measurements [30,54]. The parameter λ determines the inflaton mass at the potential minimum $V(0)$, with $m_\phi = \sqrt{2\lambda}M_P \simeq 1.6\times10^{13}$ GeV. The Hubble parameter at the horizon exit is given by $H_* = 1.5\times10^{13}$ GeV (with $\phi_* = 5.35M_P$), and at the end of inflation, $H_{\rm end} = 6.3\times10^{12}$ GeV.

To determine how gravitational particle production depends on the nonminimal coupling ξ and the reheating temperature $T_{\rm reh}$, we compute the DM relic density using Eq. (6) and impose the observational value $\Omega_{\rm DM}h^2 = 0.1198$ [35]. Since the relic abundance depends on the reheating temperature, we apply the reheating temperature limits $T_{\rm BBN} \le T_{\rm reh} \le T_{\rm max}$, where the BBN temperature is $T_{\rm BBN} \simeq 1$ MeV and $T_{\rm max} \simeq 2\times10^{15}$ GeV. We illustrate in the top panel of Fig. 2 the reheating temperature as a function of ξ for a range of masses varying from $\mathcal{O}(10)$ GeV to $\mathcal{O}(10^{12})$ GeV. We find that when $m_\chi \ll H_{\rm end}$, as $\xi \to 0$ the limit asymptotes to $T_{\rm reh} \simeq 20$ GeV. As ξ increases, the reheating temperature peaks in the range of $0.17 \lesssim \xi \lesssim 0.20$, and the conformal coupling $\xi = 1/6$ lies in this domain. We find that for $\xi > 1$, the reheating temperature function can be fitted with the expression $T_{\rm reh}(m_\chi, \xi) = a\exp(-b\xi)(1\ {\rm GeV}/m_\chi)$, where $a = 1.1\times10^{14}$ and $b = 0.24$.

Next, we use the thermal DM relic density expression (8) and display $T_{\rm reh}(m_\chi, \xi)$ in the bottom panel of Fig. 2.

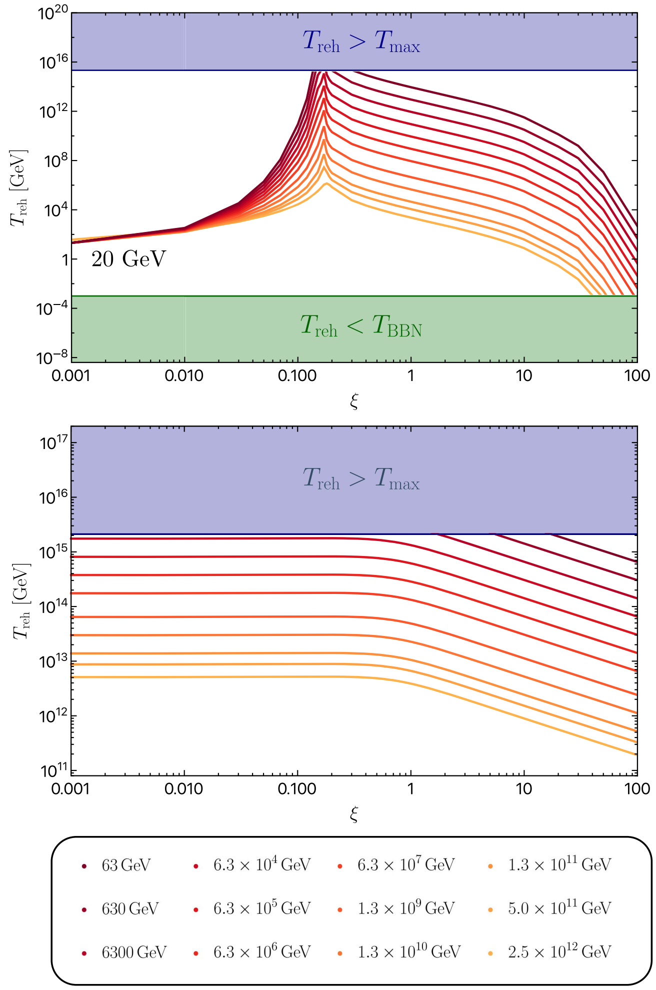

FIG. 2. Reheating temperature dependence on ξ for a range of masses. The top panel outlines gravitational production constraints, while the bottom highlights thermal production. The color code corresponds to the bare DM mass m_χ, as indicated in the legend. Regions below the respective colored lines indicate $\Omega_\chi h^2 < 0.12$.

As can be observed, the thermally produced DM requires an extremely large reheating temperature close to $T_{\rm max}$. However, when comparing it with the purely gravitational particle production (top panel), we can see that the thermal production remains subdominant throughout the entire parameter space.

To impose the isocurvature constraints, we numerically compute the isocurvature power spectrum (9) while imposing an experimental upper bound $\mathcal{P}_{\mathcal{S}}(k_*) \lesssim 8.3 \times 10^{-11}$ [35]. For minimal coupling, we find that $m_\chi \gtrsim 7.7 \times 10^{12}$ GeV $\simeq 1.2 H_{\rm end}$, which agrees with the well-known results that a scalar DM field with minimal coupling must be superheavy [16,17,21]; for a conformal coupling $\xi = 1/6$, we find that the isocurvature constraints are always satisfied, and in the massless limit we find that $\xi \gtrsim 0.02$.

Finally, we estimate the Lyman-α bound (11). We find that the DM bound increases as a function of ξ, with

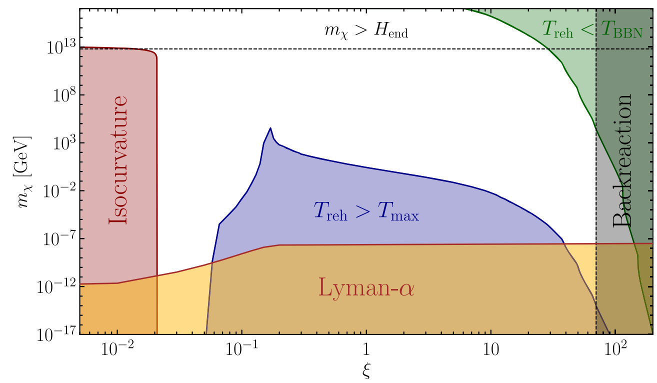

FIG. 3. Parameter space of the DM mass m_χ as a function of the nonminimal coupling ξ. The white region represents the space compatible with the observed DM abundance. The colored region shows the excluded constraints, while the gray region indicates when the backreaction effects become significant.

$m_\chi \gtrsim 2 \times 10^{-4}$ eV when $\xi = 0$, and it plateaus around $\xi \gtrsim 0.2$, with the bound given by $m_\chi > 34$ eV. The details of our derivation can be found in the Supplementary Material sections.

We show our combined parameter space in Fig. 3. In general, the nonminimal coupling ξ is constrained to $0.1 < \xi < \mathcal{O}(100)$, with the lower bound resulting from the isocurvature constraint and the upper value from the BBN constraint. We note that we have not studied the superheavy region $m_\chi > H_{\rm end}$, which would also have constraints arising from $T_{\rm BBN}$ and $T_{\rm max}$, and we plan to investigate this in future work. We found that Lyman-α constraints give the lower bound $m_\chi > 2 \times 10^{-4} - 34$ eV. The strongest constraint arises from $T_{\rm max}$ in the range of $0.1 \lesssim \xi \lesssim 1$.

In this Letter, we have explored a simple and compelling scenario involving a spectator DM scalar field that couples nonminimally to gravity. Our study demonstrates that the presence of nonminimal coupling opens up a broad parameter space, subject to constrains arising from maximum reheating temperature, BBN temperature, isocurvature, Lyman-α limits. We look forward to forthcoming experiments such as the Simons Observatory [55], CMB-S4 [56], and LiteBIRD [57], which could potentially detect B-modes in the CMB, provide more comprehensive scalar power spectrum analysis, and either strengthen isocurvature mode constraints or detect it. Since in this scenario, the DM field χ contains a blue-tilted isocurvature component, it predicts an enhancement in the power spectrum at small scales, making it amenable for being further constrained by improved structure formation and spectral distortion data [58–60].

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