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Dark matter decaying into a Fermi sea of neutrinos

Ole Eggers Bjælde[1,2,*] and Subinoy Das[3,†]

[1]Institut für Theoretische Teilchenphysik und Kosmologie RWTH, Aachen University, D-52056 Aachen, Germany
[2]Department of Physics and Astronomy, Aarhus University, Ny Munkegade 120, DK-8000 Aarhus C, Denmark
[3]Department of Physics and Astronomy, University of British Columbia, BC, V6T 1Z1 Canada
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We study the possible decay of a coherently oscillating scalar field, interpreted as dark matter, into light fermions. Specifically, we consider a scalar field with sub-eV mass decaying into a Fermi sea of neutrinos. We recognize the similarity between our scenario and inflationary preheating where a coherently oscillating scalar field decays into standard model particles. Like the case of fermionic preheating, we find that Pauli blocking controls the dark matter decay into the neutrino sea. The radius of the Fermi sphere depends on the expansion of the universe leading to a time varying equation of state of dark matter. This makes the scenario very rich and we show that the decay rate might be different at different cosmological epochs. We categorize this in two interesting regimes and then study the cosmological perturbations to find the impact on structure formation. We find that the decay may help in alleviating some of the standard problems related to cold dark matter.

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I. INTRODUCTION

Dark matter has become an extremely interesting area of research in both cosmology and particle physics. From the particle physics point of view it can be thermal WIMPs, axions, Kaluza Klein states, etc. Though supersymmetric (SUSY) models predict its mass to be of the order the electroweak scale, there are viable models of dark matter where its mass can be as low as sub-eV, for example, axionlike dark matter. Especially the direct and indirect search for dark matter has narrowed down the parameter space of these well studied candidates to a large extent. Hence it is highly likely that dark matter is of much more exotic nature than thought of. In addition, there may be a requirement for more complicated physics such as interactions with dark energy or with neutrinos. Similarity between the neutrino mass and the present dark energy density scale has inspired people to look for a connection between the two. Now we know that the normal active neutrino cannot be a viable dark matter candidate because of its free-streaming ability. But the existence of sub-eV neutrino mass might point towards richer sub-eV scale physics. In fact there have been a few interesting works [1,2] where new states of sub-eV masses are present in the dark sector. The reason is, if the dark sector interacts only gravitationally with the standard model sector, a TeV scale SUSY breaking in SM would predict scalar particles of mass $\frac{\mathrm{TeV}^2}{M_{\mathrm{Pl}}} \sim 10^{-3}$ eV in the dark sector. Scalar dark matter of milli eV mass with a possible coupling to neutrinos has been discussed in [3]. Also moduli of sub-eV mass can easily arise from string compactification [4].

If in nature dark matter arises from such a low energy scale, we would expect it to decay into light fermions like neutrinos through a Yukawa type of coupling. GeV scale dark matter decay (and annihilation) to neutrinos has drawn lots of recent interest [5–11], especially in the context of recent cosmic ray measurements and the DAMA/LIBRA experiment [12,13]. But in our case, dark matter is of sub-eV mass and the signature of its decay into neutrinos is mainly cosmological, especially in structure formation. Recently, in a different context, *the nonthermal wimp miracle* [14] was introduced where a scalar of TeV mass decays into a stable dark matter particle. So the decay of a scalar particle can lead to very rich phenomenology in cosmological context.

The possible decay of scalar dark matter into light neutrinos is an effect which could potentially help to understand the apparent surplus of power on small scales in simulations containing normal CDM [15–17]. This surplus could, for instance, be an indication that CDM is simply clustering too much on small scales and we need some mechanism to reduce their gravitational interaction. This is where the decay could play a role—see [18] for a similar idea. In this paper we study the nature of the decay and its possible signature in structure formation. As we consider an axionlike scalar dark matter, it is undergoing a coherent oscillation and might decay into neutrinos through a parametric excitation. In that case, the process will have many similarities with inflationary fermionic preheating [19–22]. Therefore, we dub the process *preheating dark matter.*

Cosmology with decaying/interacting dark matter has been an interesting topic of research [23–29] in recent times as it gives a probe to detect dark matter indirectly through its effects on structure formation. On top of that, if it decays into dark energy it might give rise to a unified

*bjaelde@physik.rwth-aachen.de
†subinoy@phas.ubc.ca

description of dark matter and dark energy. In most of the studies [30,31], the rate of energy transfer from dark matter to other components (e.g., dark energy, radiation, or neutrinos) has been empirically assumed driven by mathematical simplicity. Here we present a concrete model of dark matter decay to light fermionic states like neutrinos and then study its imprint on structure formation. In particular, we derive the decay rate of dark matter as a function of redshift using the theory of fermionic preheating. We find that in the initial stage the decay rate is faster and determined by parametric resonance. But at late times the parametric production ceases and redshifts of fermionic modes control the decay rate. This time, varying decay rate makes the phenomenology rich and offers a prominent imprint on the structure formation, something which might be experimentally probed by near future experiments.

The plan of the paper is as follows: In Sec. II we discuss the particle physics aspect of our scenario. In Sec. III, we incorporate the idea of inflationary preheating for scalar dark matter decaying into neutrinos. In Sec. IV, we identify different epochs of decay and derive the background evolution, i.e. how dark matter and neutrino energy densities evolve in presence of the decay. In Sec. V, we derive the perturbation equations for our scenario, and in Sec. VI we obtain temperature anisotropy and matter power spectra. Finally, we summarize our work in Sec. VII.

II. PHENOMENOLOGICAL MODEL

Here we discuss how our scenario fits into a particle physics setup. Scalar fields of sub-eV masses are common in different particle physics models. In many models of TeV scale gauge mediated supersymmetry breaking, a gravitationally coupled dark sector contains a scalar of sub-eV mass. Also, string compactifications generically predict axionlike scalars such as the dilaton and large numbers of moduli [32,33] whose mass can easily be in the sub-eV range.

As we are interested in a sub-eV scalar field which couples to light fermions like the standard model neutrino, our set up is inspired by the models of mass varying neutrino dark energy [34] where a sub-eV mass scalar couples to the standard model neutrino—see also [35–39]. Especially, we refer to a model of supersymmetric neutrino dark energy, where multiple scalars of such low mass are present and can couple to neutrinos. In such a theory, it has been shown that the scalar potential takes the form of the well-known hybrid inflation potential. We refer readers to [1] for details and will briefly discuss here. The Lagrangian for mass varying neutrino dark energy is given by

$$\mathcal{L} \supset m_D \nu N + \kappa ANN + \text{H.c.} + V(A), \tag{1}$$

where m_D is the Dirac mass and ν is the left chiral Weyl field representing active neutrinos. N is the right-handed heavy fermion and A is the scalar. κ is some Yukawa coupling.

In supersymmetric models of neutrino dark energy where ν, N, A are promoted to superfields l, n, a, the superpotential takes the form

$$W = \kappa ann + m_D ln. \tag{2}$$

After taking quantum corrections into account, it has been shown that this leads to a hybrid inflation kind of potential. Depending on the temperature of the universe, the scalar either remains trapped at a metastable minimum playing the role of dark energy or it rolls off and starts oscillating coherently, behaving like cold dark matter. Following a simple model [3], the Lagrangian looks like

$$\mathcal{L} = \lambda n_2 \psi_3^2 + 2\lambda n_2 \psi_2 \psi_3 + m_3 \psi_3 \nu_3 + m_2 \psi_2 \nu_2 + V_{\text{susy}} + V_{\text{soft}} + V_\epsilon + \text{H.c.}, \tag{3}$$

where

$$V_{\text{susy}} = 4\lambda^2 n_2^2 n_3^2 + \lambda^2 n_3^4, \tag{4}$$

and

$$V_{\text{soft}} = \tilde{m}_2^2 n_2^2 - \tilde{m}_3^2 n_3^2 + \tilde{a}_3 n_3^3. \tag{5}$$

The terms in V_ϵ are included in order to generate a Majorana mass for the neutrino in the vacuum. In this kind of theory the superpartner sneutrinos (here denoted by n_2, n_3) can easily be of sub-eV mass and play the role of dark matter. Also it can easily couple to light fermions (neutrinos). From now on, we will switch our focus to cosmological effects of such a model.

III. PREHEATING FROM SCALAR DARK MATTER

We are essentially interested in a light scalar field dark matter of sub eV mass which has a coupling to an ultra light fermion which for our case is the standard model neutrino. We consider decays of such dark matter into neutrinos, though our framework is true for decay into any fermion. Following the mechanism of inflationary preheating, in this section we will understand the physical nature of the decay and will clarify different regimes of decay. We will see that the decay rate changes as the universe expands due to time evolution of a resonance parameter which controls the parametric excitation of neutrinos.

We mainly follow the work [22] on fermionic preheating and apply that to our scenario. So we refer to this work for detailed derivation of the equations. Briefly, to find the number density of created fermions through the preheating mechanism in an expanding background, one derives a mode equation using the original Dirac equation with the Friedmann-Lemaître-Robertson-Walker metric. It has been shown that the comoving number density of created fermions can be obtained by solving for a mode function $X_k(t)$. For a Yukawa type coupling $\lambda \phi \psi \bar{\psi}$, the mode equation is given by

$$X_k'' + [\kappa^2 + (\tilde{m} + \sqrt{q}f)^2 - i\sqrt{q}f']X_k = 0, \qquad (6)$$

where $\phi_0 f(t)$ is the background solution for the time evolution of the oscillating scalar field, $\kappa \equiv \frac{k}{m_\phi}$ is the dimensionless fermion mass, $\tilde{m} \equiv \frac{m_\psi}{m_\phi}$, and the resonance parameter $q \equiv \frac{\lambda^2 \phi_0^2}{m_\phi^2}$. These three parameters completely determine the parametric production of fermions. We consider the oscillation of the field with the usual quadratic potential $V = \frac{1}{2}m^2\phi^2$ as this is a good approximation around minima of any potential. The term $(\tilde{m} + \sqrt{q}f)$ can be thought of as an effective mass of the fermion. As the scalar field oscillates, the effective mass itself will oscillate around zero and the parametric production of fermions is enhanced when the effective mass crosses zero. It has been shown numerically that $n_k(t)$ oscillates and, due to Pauli blocking, its maximum value never crosses unity. But for decay into bosonic particles it is not bounded by unity. We stress that the behavior of this parametric production is considerably different than the perturbative decay process $\phi \longrightarrow \bar{\psi}\psi$, where the decay rate is given by $\Gamma \simeq \frac{\lambda^2 m}{8\pi}$.

In the above mode equation, expansion of the universe has been neglected which may only be true at very late times where the Hubble parameter drops. To get a full understanding, one must include the expansion of the universe. This alters two aspects. The parameters q and κ now become time dependent. More specifically, we get $q \equiv \frac{\lambda^2\phi^2(\tau)}{m_\phi^2}$ and the physical momentum $p \equiv \frac{\kappa}{a(t)m_\phi}$ where $a(t)$ is the scale factor of the universe. As a result, the periodic modulation of the comoving number density does not hold anymore. For large values of the resonance parameter ($q \geq 1$), the calculation of parametric production becomes, in fact, simple. Luckily, we will see later that for our case $q \gg 1$ for large periods of (cosmological) time. Using the method of successive scattering for fermions, it has been shown [22] that due to the loss of periodicity of n_k the production of fermions happens through a stochastic filling of a Fermi sphere up to a Fermi radius κ_F which depends on scale factor $a(t)$ and is given by

$$\kappa_F^2 \simeq \sqrt{q(t)^{1/2}}a(t). \qquad (7)$$

Now to find the exact dependence, one needs to know how $q(t)$ changes with scale factor. As we are using a quadratic potential for the scalar, the solution for the scalar field for this case is well known. Oscillation of ϕ in this case is given by the asymptotic solution

$$\phi(t) \sim \frac{\phi_0}{a^{3/2}}\cos(t). \qquad (8)$$

Using this, it is easy to derive

$$\kappa_F = m_\phi q_0^{1/4} a^{1/4}, \qquad (9)$$

where $q_0 \equiv \frac{\lambda^2\phi_0^2}{m_\phi^2}$.

So, as the universe expands, the Fermi sphere also expands producing more and more neutrinos. But the resonance parameter decreases as the amplitude of oscillation drops due to Hubble friction and at some point the Fermi sphere stops expanding when $q(t)$ becomes of the order of unity. At this regime, the redshift of fermionic modes due to Hubble expansion is fast enough to prevent the parametric excitation. Finally, fermions will be produced with a much lower rate in the perturbative regime and perturbative processes continue unless $m_\phi < 2m_\psi$.

IV. DIFFERENT DECAY REGIMES

Using the above results, now we can focus on production of neutrinos and its time evolution. Here we assume that parametrically produced neutrinos mix with other relic neutrinos and acquire the same temperature through thermalization. It is instructive to note that as the mass of the scalar is way less than in the usual inflationary preheating scenario, we would get parametric excitation until very late times. From the previous discussion, we have learned that the resonance parameter q is very crucial to determine the nature of the decay and q itself is time dependent. Now we will discuss the two different decay regimes and the transition time between them for our simple model with a quadratic potential.

A. Regime I: Expanding Fermi radius ($q \gg 1$)

During early stages of parametric production

$$q(t) \equiv \frac{\lambda^2\phi(t)^2}{m_\phi^2} = 2\lambda^2\frac{\rho_{\rm DM}}{m_\phi^4}. \qquad (10)$$

As we are interested in scalar mass of the order of $m_\phi \sim 10^{-3}$ eV, almost all over the cosmic history until today, $\rho_{\rm DM} \geq (10^{-3}\text{ eV})^4$. Now if the coupling constant λ is of the order of unity, we still get parametric production at very late times. But for smaller couplings parametric excitation stops at earlier redshift when $q(t) \simeq 1$. Later we will take different choices of the coupling λ and study how it affects the formation of structure. The produced neutrino number density is obtained through the volume of the Fermi sphere with radius

$$\kappa_F^{\rm phys} = q^{1/4}a^{1/4} \times a^{-1} = q^{1/4}a^{-3/4}, \qquad (11)$$

where $q = \frac{\lambda^2\phi^2}{m_\phi^2}$. This can be used to calculate the neutrino density

$$\rho_\nu \simeq \int_0^{\kappa_F} d^3k = 8\pi\lambda^2\frac{1}{2}m_\phi^2\phi^2 = 8\pi\lambda^2\rho_{\rm DM}, \qquad (12)$$

where we have used $\rho_{\rm DM} = \frac{1}{2}m_\phi^2\phi^2$. It is important to note that the neutrino energy density is proportional to the local dark matter density. Using this fact and the continuity equation for the total dark matter and neutrino fluid, we can find the evolution of the dark matter energy density and hence neutrino energy density. The continuity equation for the dark matter and neutrino as a whole reads

$$\dot{\rho}_{\rm tot} + 3H\rho_{\rm tot}(1 + w_{\rm tot}) = 0, \tag{13}$$

where $w_{\rm tot} = \frac{P_{\rm tot}}{\rho_{\rm tot}}$, $\rho_{\rm tot} = \rho_\nu + \rho_{\rm DM}$ and $P_{\rm tot} = P_\nu + P_{\rm DM}$. Equation (13) can be split up into the two components

$$\dot{\rho}_{\rm DM} + 3H\rho_{\rm DM} = -Q \tag{14}$$

and

$$\dot{\rho}_\nu + 4H\rho_\nu = Q, \tag{15}$$

where Q represents the decay rate from dark matter to neutrinos and we have taken advantage of the fact that $P_{\rm DM} = 0$ and $P_\nu = \frac{1}{3}\rho_\nu$. Combining Eqs. (12), (14), and (15) we get the relations

$$\rho_{\rm DM} = \rho^i_{\rm DM}\left(\frac{a}{a^i}\right)^{-\iota} \qquad \rho_\nu = 8\pi\lambda^2\rho_{\rm DM}, \tag{16}$$

where the i denotes the value at some fixed time (e.g. today) and $\iota = \frac{3+32\pi\lambda^2}{1+8\pi\lambda^2}$. So here we clearly see that due to parametric production, dark matter no longer redshifts as $1/a^3$. Its effective equation of state changes from zero to slightly higher values. The higher the coupling λ, the higher the deviation. This effective equation of state, which corresponds to the value we would get if we did not know about the coupling between dark matter and neutrinos, can be calculated from a revised version of Eq. (14) $\dot{\rho}_{\rm DM} + 3H\rho_{\rm DM}(1 + w_{\rm eff}) = 0$. The result is

$$w_{\rm eff} = \frac{\iota}{3} - 1. \tag{17}$$

We note that for $\lambda \to 0$, it gives the right limit for the equation of state $w_{\rm eff} \to 0$.

B. Regime II: Fermi radius stops expanding ($q \simeq 1$)

In the second regime, parametric excitations weaken due to the drop in resonance parameter q. During this regime, $q \sim 1$ and the radius of the physical Fermi sphere has approached the constant value $k_F \sim m_\phi$. This means the decay is controlled by the redshifts of Fermi momentum due to the expansion. As the universe expands, the Fermi momentum drops, opening up space in the Fermi sphere. This space is immediately filled up by the scalar field decaying into neutrinos. The regime may be important for structure formation if the decay into neutrinos can cause a substantial decrease in the dark matter density. This is possible when $\rho_{\rm DM} \sim m_\phi^4$, because, in this case, decay of each DM particle to a neutrino causes a significant decrease in the dark matter energy density. As the dark matter mass is of the order sub eV in our model, this can happen only at late times. This can enhance the late ISW effect thus modifying the structure formation on large scales. Again, using Eq. (12) we easily obtain $\rho_\nu = 4\pi m_\phi^4$. From the continuity equations Eqs. (14) and (15) we get the relations

$$\rho_{\rm DM} = -\frac{\gamma}{3} + \left(\rho^i_{\rm DM} + \frac{\gamma}{3}\right)\left(\frac{a^i}{a}\right)^3 \qquad \rho_\nu = \frac{\gamma}{4}, \tag{18}$$

where $\gamma = 16\pi m_\phi^4$. So we see that the neutrino energy density is constant in this regime, only the dark matter density dilutes. This is what we expect, because if there were no decay into neutrinos, its density would simply redshift like the standard model neutrino governed by the Hubble expansion. But here, as soon as phase space opens in the Fermi sea of neutrinos due to cooling of the universe, it gets refilled by the decay from dark matter thus keeping its density constant.

In this regime the effective equation of state for dark matter can be calculated to be

$$w_{\rm eff} = \frac{\frac{\gamma}{3}}{\rho_{\rm DM}}, \tag{19}$$

which, depending on model parameters, can deviate an appreciable amount toward the present—see Fig. 1.

C. Transition redshift

In this subsection we find the transition redshift between the two epochs in terms of model parameters m_ϕ and λ. Patching the two regimes together at a scale factor a_T we arrive at the relations

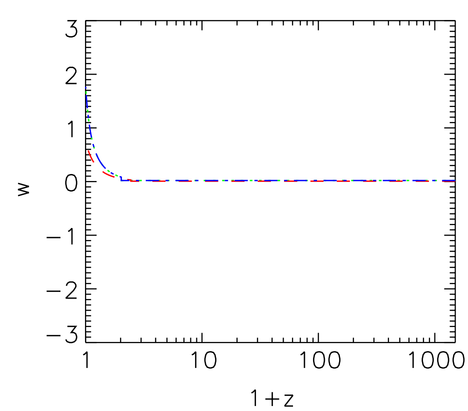

FIG. 1 (color online). The effective equation of state of dark matter in the preheating dark matter scenario with the following parameter choices: Dotted green line: $\Omega_{\rm DM,0} = 0.21$, $\lambda = 0.001$ with $\Omega_{\nu,0} = 0.022$, dot-dashed blue line: $\Omega_{\rm DM,0} = 0.06$, $\lambda = 0.05$ with $\Omega_{\nu,0} = 0.081$, and dashed red line: $\Omega_{\rm DM,0} = 0.12$, $\lambda = 0.03$ with $\Omega_{\nu,0} = 0.068$.

$$\begin{aligned}\rho_\nu &= 8\pi\lambda^2\rho_{\rm DM}\\ \rho_{\rm DM} &= \frac{m_\phi^4}{2\lambda^2}\left(\frac{a}{a_T}\right)^{-\iota} \quad \text{for } a < a_T\\ \rho_\nu &= \frac{\gamma}{4}\\ \rho_{\rm DM} &= -\frac{\gamma}{3} + \left(\rho^0_{\rm DM} + \frac{\gamma}{3}\right)\left(\frac{a_0}{a}\right)^3, \quad \text{for } a > a_T\end{aligned} \tag{20}$$

where the 0 denotes present day values, and a_T can be determined from the preheating dark matter parameters as

$$a_T = \left[\frac{(\rho^0_{\rm DM} + \frac{\gamma}{3})}{\frac{m_\phi^4}{2\lambda^2} + \frac{\gamma}{3}}\right]^{1/3} a_0. \tag{21}$$

V. PERTURBATION ANALYSIS

In order to study the implications of preheating dark matter, we perform a cosmological perturbation analysis in the synchronous gauge in which a line element is given by

$$ds^2 = a^2(\tau)[-d\tau^2 + (\eta_{ij} + h_{ij})dx^i dx^j], \tag{22}$$

where $i, j = 1, 2, 3$, η_{ij} is the Minkowski space metric, h_{ij} is the perturbation to the metric and we are using comoving coordinates $x^\mu = (\vec{x}, \tau)$ in a spatially flat background space-time. We follow the procedure given by Ref. [40], in which, to linear order in the perturbations, the stress-energy tensor is given by

$$\begin{aligned}T^0_0 &= -(\bar\rho + \delta\rho) \qquad T^0_i = (\bar\rho + \bar P)v_i = -T^i_0\\ T^i_j &= (\bar P + \delta P)\delta^i_j + \Sigma^i_j, \qquad \Sigma^i_i = 0,\end{aligned} \tag{23}$$

where the perturbations to energy density and pressure are defined as $\delta\rho = \rho - \bar\rho$ and $\delta P = P - \bar P$, Σ^i_j is the anisotropic shear perturbation, and v_i is the coordinate velocity of the fluid.[1] The latter is a small quantity and can be treated as a perturbation of the same order as $\delta\rho$ and δP. Instead of working with the velocity itself we use the divergence defined as $\theta = ik^i v_i$. Similarly, instead of the anisotropic shear perturbation, we use the shear stress σ. This is defined as $\sigma = \frac{-(k_i k^j - \frac{1}{3}\eta^j_i)\Sigma^i_j}{(\bar\rho + \bar P)}$.

The conservation of energy and momentum for our coupled fluid implies that the covariant derivative of the stress-energy tensor is 0.

$$\Gamma^{\mu\sigma}_{;\mu\,\rm fluid} = \partial_\mu T^{\mu\sigma} + \Gamma^\sigma_{\alpha\beta}T^{\alpha\beta} + \Gamma^\alpha_{\alpha\beta}T^{\sigma\beta} = 0. \tag{24}$$

However, for the individual components in the fluid it is slightly different

$$\Gamma^{\mu\sigma}_{;\mu\,\rm CDM} = -\delta Q \qquad \Gamma^{\mu\sigma}_{;\mu\ \nu} = \delta Q, \tag{25}$$

where the δQ can be determined directly from Eq. (14). Using the time-time (00) component of the stress-energy tensor from Eq. (23) and combining with Eq. (25) above, we get the equation of motion for the individual density contrasts $\delta_i = \frac{\delta\rho_i}{\rho_i}$. Using the space-space (ii) components in the same way will give us the time evolution of θ_i.

Hence we arrive at the equations of motion for the DM and neutrino components for $q \gg 1$

$$\begin{aligned}\dot\delta_\nu &= -\frac{4}{3}\left(\theta_\nu + \frac{\dot h}{2}\right)\\ \dot\delta_{\rm CDM} &= -\frac{\dot h}{2}\\ \dot\theta_\nu &= -\frac{1}{1 + 8\pi\lambda^2}H\theta_\nu + k^2\left(\frac{1}{4}\delta_\nu - \sigma_\nu\right)\\ \dot\theta_{\rm CDM} &= 0.\end{aligned} \tag{26}$$

Similarly, for $q \sim 1$ we get

$$\begin{aligned}\dot\delta_\nu &= -\frac{4}{3}\left(\theta_\nu + \frac{\dot h}{2}\right)\\ \dot\delta_{\rm CDM} &= -\frac{\dot h}{2} + 4H\frac{\rho_\nu}{\rho_{\rm CDM}}(\delta_{\rm CDM} - \delta_\nu)\\ \dot\theta_\nu &= -4H\theta_\nu + k^2\left(\frac{1}{4}\delta_\nu - \sigma_\nu\right)\\ \dot\theta_{\rm CDM} &= 0.\end{aligned} \tag{27}$$

It is the effect of the term containing λ and the term containing the neutrino energy density ρ_ν in the evolution of the CDM density contrast that separates the evolution of perturbations in the preheating dark matter case from the normal case, where the decay of CDM is not permitted.

VI. RESULTS

In this section we present the numerical results we obtained by using the equations from the two previous sections. We modified the publicly available CMBFAST code [42] to include preheating dark matter to get both the matter power spectrum and the temperature anisotropy spectrum. This code is developed to calculate the linear CMB anisotropy spectra based on integration over the sources along the photon past light cone, but also outputs transfer functions from which the linear matter power spectrum can be calculated.

In our analysis, we keep the epoch of matter-radiation equality fixed, and the only free parameters are the current value of $\Omega_{\rm DM,0}$ and the parameter λ. In addition, we keep the amount of baryons today fixed at $\Omega_b = 0.05$ and choose the normalized Hubble expansion rate at the value $h_{\rm Hubble} = 0.7$. We include one species of massless neutrinos produced in the decay as well as the three standard model neutrinos, which, for simplicity, are assumed to have a degenerate mass spectrum with $m_\nu = 1.5 \times 10^{-3}$ eV. We assume that the neutrinos produced in the decay mix with standard neutrinos, although this assumption makes no qualitative difference to the results.

[1] For further information about cosmological perturbation theory see Ref. [41]

The impact of preheating dark matter on the temperature anisotropy spectrum can be seen in Fig. 2. The most apparent difference from the spectrum of ΛCDM can be seen on largest angular scales, $l \lesssim 100$ (corresponding roughly to a degree), although for some choices of parameters the positions and relative heights of the peaks are also affected. We generally observe an increase in power on scales, $20 < l < 100$, whereas on scales $l \lesssim 10$ we see an increase or a decrease in power depending on the model parameters. For scales $l \lesssim 100$, the dominant contribution to the CMB temperature anisotropy spectrum comes from the Integrated Sachs-Wolfe effect, which arises because of the evolution of the gravitational potentials encountered by the photon on its journey from the last-scattering surface. The modification to the cosmological background because of preheating dark matter can be quite significant, as we saw in the last section. This is particularly so in the $q \sim 1$ regime, where the neutrino energy density becomes constant, which leads to a second term in the evolution of the CDM perturbation in Eq. (27).

On the largest angular scales $l < 20$, the dominant contribution to the CMB temperature anisotropy spectrum comes from the late-time Integrated Sachs-Wolfe effect (ISW). This effect is a result of the universe entering an epoch of rapid expansion as it becomes dominated by dark energy. In this epoch, photons moving into gravitational potential wells will get a boost as the potential well is decaying and becomes slightly shallower while the photon is passing through it—and vice versa for gravitational hills. It is clear that this effect depends intimately on many different parameters such as $\Omega_{\mathrm{DM},0}$ and $\Omega_{\nu,0}$. In our model, dark matter is being transformed into light neutrinos—i.e. radiation—most efficiently at late times, and hence we expect an effect on the largest scales. In the context of interacting dark matter–dark energy models, the ISW effect has been studied in [43].

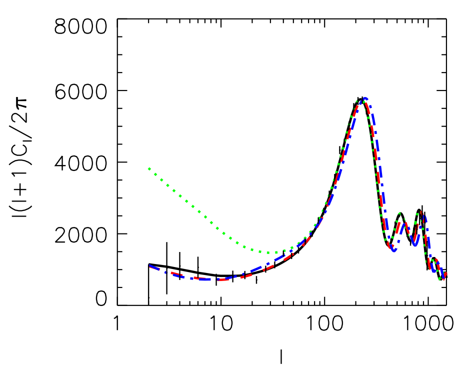

FIG. 2 (color online). The temperature anisotropy spectrum as a function of the angular modes in the preheating dark matter scenario with the following parameter choices: Solid black line: ΛCDM with $\Omega_{\mathrm{DM},0} = 0.24$, $\Omega_{\Lambda,0} = 0.71$, dotted green line: $\Omega_{\mathrm{DM},0} = 0.21$, $\lambda = 0.001$ with $\Omega_{\nu,0} = 0.022$, dot-dashed blue line: $\Omega_{\mathrm{DM},0} = 0.06$, $\lambda = 0.05$ with $\Omega_{\nu,0} = 0.081$, and dashed red line: $\Omega_{\mathrm{DM},0} = 0.12$, $\lambda = 0.03$ with $\Omega_{\nu,0} = 0.068$.

Turning our attention to the matter power spectra, we know that the linear matter power spectrum is extremely well determined by SDSS (see Ref. [44]) and WMAP (see Ref. [45]) on intermediate and large scales. And in addition, Ly-α forest data have some constraints on the small scales—see e.g. [46]. Hence we normalize our matter power spectra such that they coincide with matter power spectra obtained from using normal ΛCDM at the largest scales (these are also the latest to have entered the horizon). The results are presented in Fig. 3 for different values of λ and $\Omega_{\mathrm{DM},0}$. A small damping on small angular scales seems to be generic, similar to standard models of CDM and hot dark matter, where a similar reduction in power is achieved. We note that in order to comply with e.g. supernova data (Ref. [47]) we cannot change $\Omega_{\mathrm{M},0}$ drastically. Of course we do turn CDM into neutrinos that redshift as radiation—only slightly faster than CDM. Hence we have some room to change $\Omega_{\mathrm{M},0}$ and still be in agreement with data.

The matter power spectra for the different parameter choices are agreeing relatively well with ΛCDM as result of the small λ-value. Still, we do notice the small reduction on the smallest scales which can be probed by CMBFAST. This is to be expected since part of the CDM responsible for the gravitational wells is decaying into neutrinos which undergo free-streaming on these small scales. This pre-

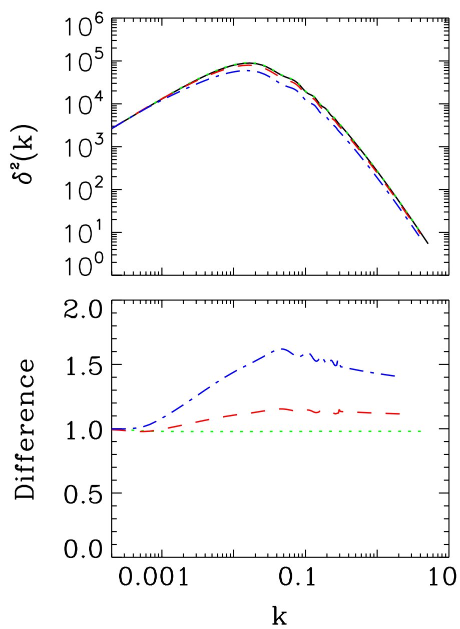

FIG. 3 (color online). Top: The matter power spectra as a function of $k[\mathrm{Mpc}^{-1}]$ in the preheating dark matter scenario with the following parameter choices: Solid black line: ΛCDM with $\Omega_{\mathrm{DM},0} = 0.24$, $\Omega_{\Lambda,0} = 0.71$, dotted green line: $\Omega_{\mathrm{DM},0} = 0.21$, $\lambda = 0.001$ with $\Omega_{\nu,0} = 0.022$, dot-dashed blue line: $\Omega_{\mathrm{DM},0} = 0.06$, $\lambda = 0.05$ with $\Omega_{\nu,0} = 0.081$, and dashed red line: $\Omega_{\mathrm{DM},0} = 0.12$, $\lambda = 0.03$ with $\Omega_{\nu,0} = 0.068$. Bottom: The difference between the preheating dark matter model and ΛCDM $[\delta^2(\Lambda\mathrm{CDM})/\delta^2(\mathrm{PHDM})]$ as a function of k.

vents them from clustering, and since we are creating an appreciable amount of neutrinos due to the decay of CDM, we generally expect this reduction of power on small scales.

On even smaller scales we also expect a considerable effect, which will reduce the clustering ability of CDM. Unfortunately, we cannot probe those scales satisfactorily with CMBFAST—as it is using linear perturbation analysis and we expect the evolution to be highly nonlinear. The investigation of the effect of preheating dark matter on nonlinear scales is beyond the scope of the present work and will be postponed to a future paper.

VII. CONCLUSION

In this work, we have considered a coherently oscillating scalar field of sub-eV mass which behaves as dark matter. Because of its coupling, it slowly decays into neutrinos as the universe expands until today. The decay rate as a function of redshift has been derived following the physics of inflationary preheating of a scalar into fermions. We find that the decay rate is modulated by Pauli blocking and the expansion of the universe, giving us rich physics of dark matter decay into a neutrino sea.

We studied the effect of the decay on structure formation and obtained spectra for the anisotropies in the cosmic microwave background and matter power spectra. For the parameters proposed in this paper, we showed that given the decay we are able to slightly reduce the amount of power on small scales in the matter power spectra—something which seems to be required from data—while at the same time being in good agreement with SDSS and WMAP observations.

Interestingly, the proposed scenario leads to features in the temperature anisotropy spectrum,which can be seen as a prominent late ISW effect and slight modifications to the second and third peaks. Consequently, as future direction, it would be interesting to study the late ISW effect in detail, from which we will be able to constrain the model parameters more effectively. In addition, we would like to do a follow-up COSMOMC analysis using the newest data available as well as to examine the effect of preheating dark matter on structure formation on nonlinear scales. After such an analysis it will be more clear what the best choice of preheating dark matter parameters is such that a comparison with future Planck data, for instance, will be easier. Furthermore, we expect future weak lensing surveys to be useful in constraining our scenario since they can provide insight into the statistics of the dark matter distribution—hence they can (hopefully) shed some light on what happens on small (and large) scales of gravitational clustering.

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