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(Standard model) universe dominated by the right matter

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We analyze the phenomenology of a prolonged early epoch of matter domination by an unstable but very long-lived massive particle. This new matter domination era can help to relax some of the requirements on the primordial inflation. Its main effect is the huge entropy production produced by the decays of such a particle that can dilute any possible unwanted relic, as the gravitino in supersymmetric models, and thus relax the constraints on the inflationary reheating temperature. A natural candidate for such a heavy, long-lived particle already present in the standard model of the electroweak interactions would be a heavy right-handed neutrino. In this case, we show that its decays can also generate the observed baryon asymmetry with right-handed neutrino masses well above the bound from gravitino overproduction.

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I. INTRODUCTION

Inflation was introduced in the 80s [1] as a solution to several problems of the big bang cosmology. Perhaps the main problem was the large-scale smoothness problem. Why do different patches of the Universe, that were not in causal contact in the radiation last scattering era, have approximately the same temperature today [2]? In inflationary models, an epoch of exponential expansion inflates a small patch of the Universe in causal contact to contain all the observable Universe today. Simultaneously if the temperature after inflation is low enough, inflation helps also to dilute unwanted relics from higher scales and reduces the flatness problem.

It is usually assumed that some kind of inflation starts already at the Planck scale to avoid the Universe collapse in a few Planck times if $\Omega > 1$ or (for any Ω) to prevent the invasion of the surrounding inhomogeneity to our homogeneous patch before inflation. On the other hand, the scales observable today in the cosmic microwave background left the horizon at an energy $V^{1/4} \lesssim 6 \times 10^{16}$ GeV [3], or 60 e-foldings before the end of inflation. So that, inflation must end below this scale. After the end of inflation comes an era of reheating when the inflaton field oscillates around its minimum and decays to ordinary particles. The final reheating temperature where we recover ordinary big bang cosmology can take any value from $V^{1/4}$ above to scales as low as 1 MeV. However, the required dilution of unwanted relics, as grand unified theory (GUT) monopoles or gravitinos in supersymmetric models, forces the reheating temperature to be well below the GUT scale or even below $T_{RH} \leq 10^8$ GeV in supersymmetry models.

In this paper, we propose a simple and economical mechanism that helps solve some of these problems and reduces the requirements on the primordial inflationary mechanism without further additions to the particle spectrum of the standard model *with right-handed neutrinos*.[1] After an initial inflationary epoch (still necessary to reproduce the observed correlation on temperature fluctuations at large scales) we assume our Universe is radiation dominated for a short period and then enters a matter domination era due to the existence of a heavy long-lived unstable particle that decays to radiation well-before nucleosynthesis, when we connect with usual cosmology. In the standard model, as we will show, this role could be played by a heavy right-handed neutrino. In the literature, it is well-known that late time entropy release can help to ameliorate some of the problems of standard cosmology. However, so far most of these works have only considered moduli fields in supersymmetric theories (see for instance [4–8]) and their real presence in nature could be considered more speculative than the existence of right-handed neutrinos.

As we show below, this matter domination mechanism, naturally embedded in the standard model (SM), is able to help primordial inflation in several aspects. A long period of matter domination can reduce mildly the number of e-folds before the end of inflation at which observable perturbations were generated, relaxing this way flatness conditions on the inflationary potential. Moreover, the large entropy production in the decay of this particle completely dilutes any unwanted relics, eliminating the constraint on the inflation reheating temperature. In this sense our matter domination epoch has the same advantages as thermal inflation [9], without resorting to yet another scalar field and/or scalar potential.

[1] From now on, we call the standard model with right-handed neutrinos simply ‘‘standard model’’

II. REWRITING THE HISTORY OF THE UNIVERSE

Let us assume for a moment that at an early time a massive particle dominated the energy balance of the Universe by many orders of magnitude. How does the observed Universe feel this new epoch? Are there any observable consequences of this?

As it is well-known [10,11], a massive particle X becomes nonrelativistic when the temperature of the thermal bath falls bellow its mass, M_X, and its energy density freezes out when it drops out of equilibrium. Then, X-relic abundance relative to photons (radiation) becomes constant and therefore, if X is completely stable, the energy density of the X particles will eventually become larger than the radiation one, dominating the energy balance of the Universe. If X is not completely stable but rather sufficiently long-lived to dominate the energy density, later on X will decay into relativistic particles that thermalize. The radiation content of the Universe will be increased and it will enter again in a radiation dominated era. This will be the basic evolution of the Universe in our model.

We start from a radiation dominated Universe at a scale larger than the mass of our X particle, M_X, where this particle is in thermal equilibrium. We assume that this particle decouples from the plasma at a temperature of the order of its mass. This particle is unstable, although very long-lived, i.e. it has very weak interactions with radiation degrees of freedom. Then, its energy density, ρ_X, starts diluting as matter, much slower than radiation. If its lifetime, Γ_X, is long enough, it will necessarily dominate the energy density of the Universe. Although our X particle decays all the time through an exponential law [12], it is only when the age of the Universe is of the order of $1/\Gamma_X$ that the decay will sizably reduce its abundance. Once it reaches this point, it will quickly decay into radiation and our Universe will go back to a radiation dominated epoch where it will connect with the usual cosmology. This "matching" with the standard scenario must happen well-before nucleosynthesis.

The evolution equations for the matter and radiation energy-densities are well-known

$$\dot{\rho}_X = -3H(1+w_X)\rho_X - \Gamma_X\rho_X, \tag{1}$$

$$\dot{\rho}_r^{\,\mathrm{old}} = -4H\rho_r^{\mathrm{old}}, \tag{2}$$

$$\dot{\rho}_r^{\,\mathrm{new}} = -4H\rho_r^{\mathrm{new}} + \Gamma_X\rho_X, \tag{3}$$

$$H^2 = \frac{\dot{a}}{a} = \frac{8\pi}{3M_{Pl}^2}(\rho_r^{\mathrm{old}} + \rho_r^{\mathrm{new}} + \rho_X), \tag{4}$$

where w_X is the equation of state parameter of particle X, which drops from $1/3$ to zero as X becomes nonrelativistic, ρ_r^{old} is the energy density in radiation not related to X decays, ρ_r^{new} is the energy density in radiation produced by X decays, and here we assume a flat Universe after inflation consistent with observations [2].

From these equations we can see that, when X is nonrelativistic, the number of X's per comoving volume ($N_X = R^3\rho_X/M_X$) follows a simple exponential decay law and the (formal) solution to these equations is given by

$$\rho_X = \rho_X^0\left[\frac{a}{a^0}\right]^{-3} e^{-\Gamma_X(t-t_0)}, \tag{5}$$

$$\rho_r^{\mathrm{old}} = \rho_r^{\mathrm{old}^0}\left[\frac{a}{a^0}\right]^{-4}, \tag{6}$$

$$\rho_r^{\mathrm{new}} = \rho_X^0\left[\frac{a}{a^0}\right]^{-4}\int_{t_0}^{t} dt'\left[\frac{a(t')}{a^0}\right] e^{-\Gamma_X t'}\Gamma_X, \tag{7}$$

$$H^2 = \frac{\dot{a}}{a} = \frac{8\pi}{3M_{Pl}^2}(\rho_r^{\mathrm{old}} + \rho_r^{\mathrm{new}} + \rho_X), \tag{8}$$

where the superscript zero denotes the value of that quantity at the initial epoch.

In general, it is not possible to integrate analytically these equations, although useful approximations exist [12,13]. However, it is always possible to solve them numerically as we have done to generate Figs. 1 and 2. These evolution equations, (1) to (4), were thoroughly analyzed by M. Turner and collaborators in Refs. [12,13] (for a more recent work see for example [14]) and we agree completely with their analysis of the matter and radiation densities, entropy, and temperature. However, we are especially interested in the particular limit of small Γ_X and large M_X. In fact, as we will see below, we will focus on the limit of small Γ_X keeping $\Gamma_X \gg 1/t_{BBN}$ so that no trace of X is present at nucleosynthesis time, in agreement with experimental observations.

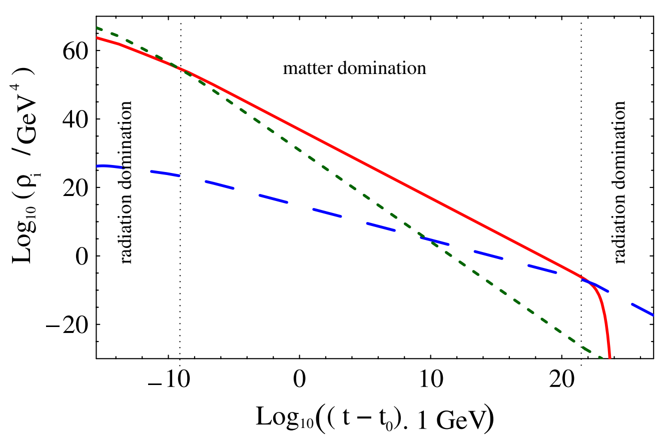

FIG. 1 (color online). Evolution of the different components of the energy density of the Universe from $T \sim 10^{14}$ GeV to $T \sim 1$ MeV. The short-dashed (green) line corresponds to ρ_r^{old} versus time, the long-dashed (blue) line to ρ_r^{new}, and the solid (red) line to ρ_X. We start from $\rho_X = \rho_r^{\mathrm{old}}/200$ at $T \sim 10^{15}$ GeV, with a $\Gamma_X = 10^{-20}$ GeV.

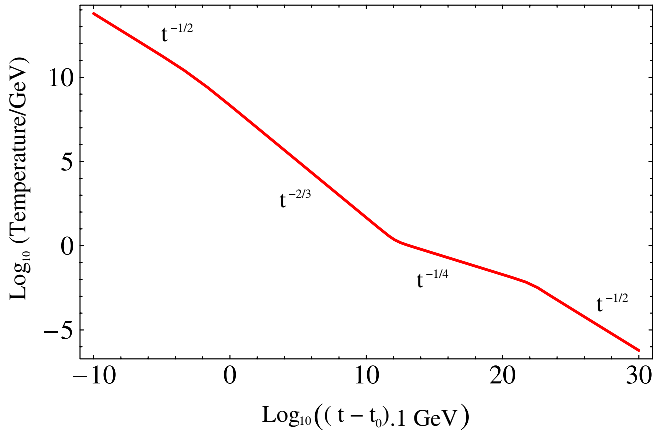

FIG. 2 (color online). Temperature of the Universe (in GeVs) as a function of time. The four different epochs in the evolution of the Universe, radiation domination, matter domination, decay, and radiation domination again can be seen in the different slopes of the curve. The time dependence of the temperature is explicitly indicated for each epoch.

In this limit, our massive particle X dominates the energy density of the Universe by many orders of magnitude during a sizeable fraction of the thermal history of the Universe (see Fig. 1). As shown in Refs. [12,13], during the X decays the temperature of the Universe does not fall as $t^{-1/2}(a^{-1})$, but rather as $t^{-1/4}(a^{-3/8})$ due to the entropy release of the decays and the temperature reaches an almost flat plateau from the point where the energy density in new radiation born through X decays and the energy density in old radiation becomes comparable up to $t \simeq \Gamma_X^{-1}$. After this time, X rapidly decays and the temperature falls again as $t^{-1/2}(a^{-1})$ (see Fig. 2).

Once it has completely decayed away, our Universe is left with a temperature

$$T_{\text{post-decay}} \simeq 1.0 \cdot 10^9 \left(\frac{g_*}{200}\right)^{-1/4} \left(\frac{\Gamma}{1\ \text{GeV}}\right)^{1/2} \text{GeV}, \quad (9)$$

where g_* counts the effective number of relativistic degrees of freedom. Notice that the temperature after X decay depends only on the decay width Γ. The ratio of entropy per comoving volume before and after X decay is given by

$$\frac{S_{\text{post-decay}}}{S_{\text{pre-decay}}} \simeq 0.14 r \left(\frac{g_*}{200}\right)^{-3/4} \left(\frac{1\ \text{GeV}}{\Gamma}\right)^{1/2} \left(\frac{M_X}{10^{10}\ \text{GeV}}\right), \quad (10)$$

where $r = g_X/2$ if X is a boson and $r = 3g_X/8$ if it is a fermion, with g_X the total number of spin degrees of freedom.

As mentioned before, an early period of matter domination, triggered by a long-lived massive particle that goes out of equilibrium and comes to dominate the energy density of the Universe before decaying, reduces the number of e-foldings before the end of inflation at which our present Hubble scale equaled the Hubble scale during inflation, i.e. the time of horizon crossing. The reduction is given by

$$\Delta N = \frac{1}{12} \ln\left(\frac{\rho_{\text{post-matter-}d}}{\rho_{\text{pre-matter-}d}}\right), \quad (11)$$

where $\rho_{\text{pre-matter-}d}$ and $\rho_{\text{post-matter-}d}$ are the energy densities at the beginning and end of the matter dominated era, respectively. The reduction can also be expressed in terms of the X parameters as

$$\Delta N \simeq -\frac{1}{6} \ln\left(90 g_*^{-3/2} r^2 \frac{M_X^2}{\Gamma M_{Pl}}\right). \quad (12)$$

This reduction is illustrated in Fig. 3, where it can be clearly seen that to determine the number of e-foldings after the horizon crossing of a given cosmological scale, as the present Hubble scale, the complete thermal history of the Universe must be used. From nucleosynthesis onwards this history is well in place. However earlier epochs are still very uncertain. The standard cosmological model assumes that inflation gives way to a long period of radiation domination (we neglect here the period of reheating that immediately follows inflation and assume sudden transitions between the different regimes). The radiation dominated epoch lasts until a redshift of a couple of thousands before entering an era of matter domination, which at redshift below one gives way to the current acceleration.

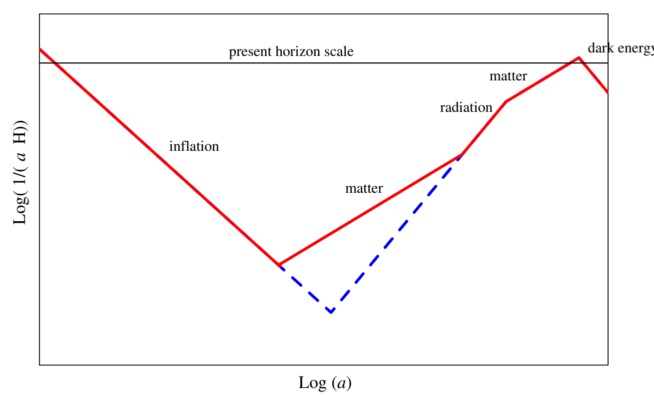

FIG. 3 (color online). Comoving Hubble radius, $\log(1/aH)$, versus $\log a$. This plot shows the different eras entering the e-foldings calculation. Inflation is an epoch where $\log(1/aH)$ is decreasing. Exponential inflation gives a line with a slope of -1. In all other cases the inflation line is shallower. During matter domination $(1/aH) \propto a^{1/2}$, while during radiation domination $(1/aH) \propto a$. The current dark energy domination signals a new inflationary epoch. The horizontal (black) solid line indicates the present horizon scale. The number of e-foldings before the end of inflation at which observable perturbations were born is the horizontal distance between the time when $(1/aH)$ first crosses that value and the end of inflation. The solid (red) line represents a Universe with a period of matter domination before big bang nucleasynthesis. The dashed (blue) line represents the standard cosmological history of the Universe with only one (recent) epoch of matter domination.

Changing the sequence of events after inflation can therefore have a strong impact on the number of e-foldings calculation. If our Universe goes through a long period of a regime where $\log(1/aH)$ scales as a^n, i.e. $H \propto a^{-(n+1)}$, it is straightforward to see that with $n > 1$ the total number of e-foldings will be increased while for $n < 1$ this number will be reduced. A period of matter domination belongs to the latter class, as in a matter dominated epoch $(1/aH) \propto a^{1/2}$ opening the door to a significant reduction on the number of e-foldings.

If we put all this information together we can see what are the required features of our particle X, its mass M_X, and its lifetime Γ, if it is to dominate over the energy density of the Universe for a long period.

Although we do want a prolonged period of matter domination, we want it to come to an end *at the latest* shortly before nucleosynthesis, as the Universe must have attained thermalized radiation domination by that time. Using Eq. (9), this condition sets a lower bound on Γ,

$$\Gamma \geq 4.6 \cdot 10^{-25}\left(\frac{g_*}{10.75}\right)^{1/2} \text{ GeV}. \tag{13}$$

We can also get an upper bound on Γ, by requesting the reheating temperature $T_{\text{post-decay}}$ to be at most 10^8 GeV, so that no unwanted relics will be produced after X decay in supersymmetric models. Such a condition reads

$$\Gamma \leq 0.7\left(\frac{g_*}{200}\right)^{1/2} \text{ GeV}. \tag{14}$$

However, for lifetimes this short (large widths), we can see from Eq. (10) that only large X masses are capable of effectively diluting the unwanted relics. In the case of the gravitino this is even more difficult as the gravitino abundance is also proportional to the temperature, $T \geq M_X$.

Of course detailed bounds can be set only after specifying the basic physics behind X, its production mechanism, its decays, or in short, its interactions. Nevertheless, we can say that requiring at least 5 orders of magnitude of dilution by entropy production for $M_X \leq 10^{10}$ GeV would need

$$2.0 \cdot 10^{-6}\left(\frac{g_*}{200}\right)^{1/2} \text{ GeV} \geq \Gamma \geq 2.0 \cdot 10^{-24}\left(\frac{g_*}{200}\right)^{1/2} \text{ GeV}. \tag{15}$$

In this case the 5 orders of magnitude of increase in the entropy should be enough to get rid of unwanted relic which could have been produced at earlier times for $T_{\text{post-decay}} \lesssim 10^8$ GeVs. Shorter lifetimes (larger widths) are also possible if we do not need such a large entropy production.

With regards to the reduction in the number of e-foldings, only a very prolonged period of matter domination, i.e. large M_X and Γ in the lower part of the allowed range, is required to give a significant reduction. In most cases, however, this number is expected to be below 10.

III. MATTER DOMINATION IN THE STANDARD MODEL

Now, we must check whether this early matter domination epoch could exist in the context of the standard model of the strong and electroweak interactions or any of its extensions. Clearly to obtain such an early matter domination era we need a massive particle, X, in thermal equilibrium at temperatures above its mass with a very long lifetime. The simplest candidate for our X particle would be a right-handed neutrino. Right-handed neutrinos are one of the minimal additions to the standard model to reproduce the observed neutrino masses through the seesaw mechanism [15]. These right-handed neutrinos, R_i, have superheavy masses, that can be as high as the grand unification scale, and are singlets under the SM gauge group. The only renormalizable couplings of R_i with the SM particles are, possibly small, Yukawa couplings. Therefore, if these couplings are sufficiently small, it seems possible that the right-handed neutrinos play the role of the X-particle with a long lifetime.

More precisely, the right-handed neutrino masses and Yukawa couplings have to reproduce the measured neutrino masses and mixings though the seesaw mechanism, $m_{\nu_L} = v_2^2 Y_\nu \cdot (M_R)^{-1} \cdot Y_\nu^T$, with v_2 the vacuum expectation value of the up-type Higgs. From here it is straightforward to obtain the required right-handed Majorana matrix from the seesaw formula itself

$$M_R = v_2^2 Y_\nu^T \cdot (m_{\nu_L})^{-1} \cdot Y_\nu. \tag{16}$$

From the light-neutrino mass matrix, m_{ν_L}, we know the mixings and the two mass differences.

The mixing matrix U is close to the so-called tribimaximal mixing

$$U = \begin{pmatrix} \frac{2}{\sqrt{6}} & \frac{1}{\sqrt{3}} & 0 \\ -\frac{1}{\sqrt{6}} & \frac{1}{\sqrt{3}} & -\frac{1}{\sqrt{2}} \\ -\frac{1}{\sqrt{6}} & \frac{1}{\sqrt{3}} & \frac{1}{\sqrt{2}} \end{pmatrix}. \tag{17}$$

Then, $m_{\nu_L} = U^* \cdot \text{Diag}(m_1, m_2, m_3) \cdot U^\dagger$, and the inverse of this matrix is $m_{\nu_L}^{-1} = U \cdot \text{Diag}(\frac{1}{m_1}, \frac{1}{m_2}, \frac{1}{m_3}) \cdot U^T$. Therefore

$$m_{\nu_L}^{-1} = \frac{1}{m_3}\begin{pmatrix} 0 & 0 & 0 \\ 0 & \frac{1}{2} & -\frac{1}{2} \\ 0 & -\frac{1}{2} & \frac{1}{2} \end{pmatrix} + \frac{1}{m_2}\begin{pmatrix} \frac{1}{3} & \frac{1}{3} & \frac{1}{3} \\ \frac{1}{3} & \frac{1}{3} & \frac{1}{3} \\ \frac{1}{3} & \frac{1}{3} & \frac{1}{3} \end{pmatrix} + \frac{1}{m_1}\begin{pmatrix} \frac{2}{3} & -\frac{1}{3} & -\frac{1}{3} \\ -\frac{1}{3} & \frac{1}{6} & \frac{1}{6} \\ -\frac{1}{3} & \frac{1}{6} & \frac{1}{6} \end{pmatrix}. \tag{18}$$

Experimentally we have that $m_3 = m_{\text{atm}} \simeq 0.05$ eV, $m_2 = m_{\text{sol}} \simeq 0.008$ eV, and $m_1 \ll m_2$ in the normal hierarchy situation and $m_2 = m_{\text{atm}} \simeq 0.05$ eV, $m_1 = m_{\text{atm}} - m_{\text{sol}}/2 \simeq 0.046$ eV, and $m_3 \ll m_2$ in the inverse hierarchy case [16].

As seen in Eq. (16) the masses of the right-handed neutrinos reproducing the observed light-neutrino masses and mixings are determined by the Yukawa matrix, Y_ν. Choosing the basis of diagonal $Y_\nu^\dagger Y_\nu$ and diagonal charged lepton Yukawa matrix, we have $Y_\nu = V_L \cdot \mathrm{Diag}(y_1, y_2, y_3)$. Obviously the physics depends strongly on the form of Y_ν, both on the eigenvalues, y_i, and the V_L matrix. Let us first analyze the role of V_L. We have two limiting situations: a) V_L has large mixings and is the source of the observed Pontecorvo-Maki-Nakagawa-Sakata matrix in neutrino mixings, $V_L \simeq U^*$ and b) the mixings in V_L are small, similarly to the situation observed in the Cabibbo-Kobayashi-Maskawa mixing matrix, $V_L \simeq \mathbb{1}$.

Case a) is very simple, the seesaw mechanism plays no role in the generation of the neutrino mixings. The observed large neutrino mixings are already present in the Yukawa couplings before the seesaw mechanism. This corresponds to the situation where, the light-neutrino Majorana mass matrix and the neutrino Yukawa couplings, or equivalently, the right-handed neutrino Majorana matrix and the Yukawa combination $Y^\dagger Y$, can be simultaneously diagonalized. So, we have $M_R = v_2^2 \mathrm{Diag}(y_1^2/m_1, y_2^2/m_2, y_3^2/m_3)$. The decay widths of the right-handed neutrinos will be given by $\Gamma_i = \frac{1}{8\pi} M_i (Y^\dagger Y)_{ii}$ and we must compare it with the Hubble rate, $H(T = M_i)$ in order to know if/when our massive neutrino will go out of equilibrium. Equivalently, we can compare the effective mass $\tilde{m}_i = (Y^\dagger Y)_{ii} v^2/M_i$ (i.e. $\Gamma = \frac{\tilde{m}_i}{8\pi} \frac{\tilde{M}_i^2}{v_2^2}$) with the critical mass $m_* = 1 \times 10^{-3}$ eV [17,18]. A right-handed neutrino would dominate the energy density if $\tilde{m}_i < m_*/g_*^2$ where g_* is the number of radiation degrees of freedom at $T = M_i$. The presence of g_*^2 is due to the fact that as we can see from Eq. (5), the ratio of matter and radiation densities grows as a. Therefore the universe has to expand a factor of g_* since $T = M_i$ to compensate g_* radiation degrees of freedom. In radiation dominance $H \propto a^{-2}$ and hence the lifetime has to be g_*^2 times longer.[2] Altogether, in case a), it is clear $\tilde{m}_3 = m_3$, $\tilde{m}_2 = m_2$, and $\tilde{m}_1 = m_1$. Therefore, in the normal hierarchy case, taking $m_1 \leq 10^{-10}$ eV, R_1 would dominate the energy density of the Universe with a mass of

$$M_{R_1} = \left(\frac{y_1}{10^{-6}}\right)^2 \left(\frac{1 \times 10^{-10}\ \mathrm{eV}}{m_1}\right) 6 \times 10^{11}\ \mathrm{GeV}. \tag{19}$$

This case is therefore a perfect example of how a right-handed neutrino can dominate the energy density of the Universe after inflation. In terms of $\tilde{m}_i$ we can write Eq. (10) as

$$\frac{S_{\text{post-decay}}}{S_{\text{pre-decay}}} \simeq 0.4 r \left(\frac{g_*}{200}\right)^{-3/4} \left(\frac{1 \times 10^{-6}\ \mathrm{eV}}{\tilde{m}_i}\right)^{1/2}. \tag{20}$$

[2]However, as we see below, this large effective mass, $\tilde{m}_i \simeq 10^{-5}$ eV, does not generate sufficient entropy.

Therefore, if we want 2 orders of magnitude of entropy production, we would need $m_1 = 10^{-10}$ eV (corresponding to $\Gamma = 4.6 \times 10^{-2}$ GeV) and $T_{\text{post-decay}} \simeq 2 \times 10^8$ GeV. Here, as the gravitino abundance is approximately linear with the reheating temperature [19], 2 orders of magnitude of dilution by the entropy production would correspondingly relax the bound on the reheating temperature by 2 orders of magnitude. Naturally, this can be easily improved by choosing smaller m_1 and y_1 in Eqs. (19) and (20).

However this simple model has several phenomenological problems. First, given that $\tilde{m}_1 \ll m_*$, this right-handed neutrino would not be produced thermally through its Yukawa interactions. In fact this will be a common problem of any massive particle dominating the energy density of the Universe as necessarily its decay/production rate will be much slower than the Hubble rate. Therefore we will always need another active interaction to produce our right-handed neutrino in the thermal plasma after inflation. This role could be played, for instance, by a gauged B-L interaction. Many grand unified models based on $SO(10)$ or groups containing it have an intermediate scale of the order of 10^{13} GeVs with a intermediate gauge group containing $U(1)_{B-L}$, as $SU(2)_L \times SU(2)_R \times SU(3)_c \times U(1)_{B-L}$ for example [20–22]. In these grand unified models the B-L coupling unifies with the other gauge couplings at M_{GUT} and therefore it is always strong enough to keep the right-handed neutrinos in thermal equilibrium in the unbroken phase. However, the B-L gauge interaction can never mediate the neutrino decay as it couples only diagonally in flavor. The neutrino decay would require a Yukawa interaction to lighter states that, as shown above, in this case is very small. A second problem, more specific of this particular case is that the decay of this right-handed neutrino erases completely any previously existing baryon or lepton asymmetry and therefore, we need some mechanism to generate the observed baryon asymmetry. The baryon asymmetry generated by this completely out-of-equilibrium decay of the right-handed neutrino is given by [23]

$$\eta_{\mathrm{B}} = \frac{8}{23} \varepsilon \simeq \frac{1}{16\pi} \sum_{j \neq 1} \frac{\mathrm{Im}[(Y^\dagger Y)_{1j}^2]}{(Y^\dagger Y)_{11}} \frac{M_1}{M_j}, \tag{21}$$

but, as $Y^\dagger Y$ is completely diagonal in the basis of diagonal right-handed neutrino masses, no new lepton asymmetry is generated by R_1 decays.

However, this situation is very unstable and a slight departure from the perfect alignment of the Majorana and Yukawa matrices changes the situation. If we call R the rotation diagonalizing the neutrino Yukawas in the basis of diagonal left-handed neutrino Majorana masses, we have $Y_\nu = U^* \cdot R \cdot \mathrm{Diag}(y_1, y_2, y_3)$ and

$$(M_R)_{ij} = v_2^2 y_i y_j \left(\frac{1}{m_1} R_{1i} R_{1j} + \frac{1}{m_2} R_{2i} R_{2j} + \frac{1}{m_3} R_{3i} R_{3j} \right). \tag{22}$$

If the Yukawa couplings are sufficiently hierarchical, $y_3 \gg y_2 \gg y_1$, the heaviest eigenvalue will be given by the element $(M_R)_{33}$,

$$(M_R)_{33} \simeq v_2^2 y_3^2 \left(\frac{1}{m_1} (\sin\theta_{13})^2 + \frac{1}{m_2} (\sin\theta_{23})^2 + \frac{1}{m_3} (\cos\theta_{13} \cos\theta_{23})^2 \right), \tag{23}$$

where we used the standard Particle Data Group parametrization for the matrix R [24]. From here, we can see that the contribution from $m_1 \ll m_2, m_3$, will dominate $(M_R)_{33}$, and hence the heaviest right-handed neutrino eigenvalue, if $(\sin\theta_{13})^2 > m_1/m_3$. For $m_1 = 1 \times 10^{-10}$ eV, a $\sin\theta_{13} > 4.4 \times 10^{-5}$ will be enough and such a small departure from perfect alignment will completely change the situation. To understand this, it is enough to analyze a simpler situation with $\theta_{12} = \theta_{23} = 0$ and $\theta_{13} \neq 0$. Let us take $y_3 \simeq 1$, $y_1 \simeq 10^{-6}$ (similar to the up-quark hierarchy), $m_3 \simeq 0.05$ eV and $m_1 \simeq 10^{-10}$ eV. In this case, the two right-handed neutrino eigenvalues (the other one is unchanged) are given by

$$M_{R_3} \simeq v_2^2 y_3^2 \left(\frac{\cos^2\theta_{13}}{m_3} + \frac{\sin^2\theta_{13}}{m_1} \right), \qquad M_{R_1} \simeq v_2^2 y_1^2 \frac{1}{m_1 \cos^2\theta_{13} + m_3 \sin^2\theta_{13}}. \tag{24}$$

If $\sin\theta_{13} \gg m_1/m_3$ then both M_3 and $\tilde{m}_3$ are fixed by m_1, while M_1 and $\tilde{m}_1$ are fixed by m_3. So, we have, for m_1 in the interesting range

$$M_{R_3} = \left(\frac{y_3}{1} \right)^2 \left(\frac{1 \times 10^{-10}\ \text{eV}}{m_1} \right) \left(\frac{\sin\theta_{13}}{0.005} \right)^2 1.5 \times 10^{19}\ \text{GeV}, \qquad \tilde{m}_3 = m_1/\sin^2\theta_{13}. \tag{25}$$

And this is the only neutrino that can dominate the energy density of the Universe. Clearly, this situation is not interesting phenomenologically as it is not possible to produce it after inflation and (probably) it does not produce a large amount of entropy.

The most interesting situation corresponds to $\sin\theta_{13} < m_1/m_3$. In this case from Eq. (24) we have $M_{R_3} \simeq v_2^2 y_3^2 / m_3 (1 + \sin^2\theta_{13} m_3/m_1)$ and $M_{R_1} \simeq v_2^2 y_1^2 / (m_1 (1 + \sin^2\theta_{13} m_3/m_1))$. Now, the rotation that diagonalizes the right-handed mass matrix is given by $\sin\phi \simeq \sin\theta_{13} (y_1 m_3)/(y_3 m_1)$. Therefore in the basis of diagonal right-handed neutrino masses we have,

$$Y^\dagger Y = \begin{pmatrix} y_1^2 (1 + \sin^2\theta_{13} (\frac{m_3}{m_1})^2) & 0 & -y_1 y_3 \frac{m_3}{m_1} \sin\theta_{13} \\ 0 & y_2^2 & 0 \\ -y_1 y_3 \frac{m_3}{m_1} \sin\theta_{13} & 0 & y_3^2 \end{pmatrix}' \tag{26}$$

and for $\sin\theta_{13} < m_1/m_3$ we have $\tilde{m}_1 \simeq m_1$ and $\tilde{m}_3 \simeq m_3$. This means the lightest right-handed neutrino, with a mass approximately given by Eq. (19), can still dominate the energy density. In such a situation, we can see from Eq. (21), that the generated baryon asymmetry is given by

$$\eta_B \simeq \frac{1}{16\pi} \frac{\text{Im}[(y_1 y_3 \frac{m_3}{m_1} \sin\theta_{13})^2]}{(y_1^2)} \frac{y_1^2 m_3}{y_3^2 m_1} \simeq \frac{1}{16\pi} y_1^2 \sin^2\theta_{13} \frac{m_3^3}{m_1^3}. \tag{27}$$

Taking $y_1 \simeq 10^{-7}$, $m_3 \simeq 0.05$ eV, $m_1 \simeq 10^{-12}$ eV, and $\sin\theta_{13} = 0.1 \times m_1/m_3$, we would obtain $\eta_B \simeq 10^{-7} \sin\varphi$ with φ the CP violating phase of $(Y^\dagger Y)_{13}^2$. Therefore, in this case, it would be possible to generate the observed baryon asymmetry and simultaneously dilute the relic density and, in particular, the gravitino density by 3 orders of magnitude. This means that the bound on the inflationary reheating temperature would be relaxed by 3 orders of magnitude. Clearly, using smaller values for m_1 and y_1 the situation can be improved arbitrarily.

The second example of Yukawa mixing matrix was case b) where $V_L \simeq \mathbb{1}$ so that we can neglect this rotation on m_{ν_L} as a small rotation will not modify the contribution of the different neutrino eigenvalues to the matrix elements. Then, we have

$$M_R = v^2 \begin{pmatrix} y_1^2 \frac{m_2^{-1} + 2m_1^{-1}}{3} & y_1 y_2 \frac{m_2^{-1} - m_1^{-1}}{3} & y_1 y_3 \frac{m_2^{-1} - m_1^{-1}}{3} \\ y_1 y_2 \frac{m_2^{-1} - m_1^{-1}}{3} & y_2^2 \frac{m_1^{-1} + 2m_2^{-1} + 3m_3^{-1}}{6} & y_2 y_3 \frac{m_1^{-1} + 2m_2^{-1} - 3m_3^{-1}}{6} \\ y_1 y_3 \frac{m_2^{-1} - m_1^{-1}}{3} & y_2 y_3 \frac{m_1^{-1} + 2m_2^{-1} - 3m_3^{-1}}{6} & y_3^2 \frac{m_1^{-1} + 2m_2^{-1} + 3m_3^{-1}}{6} \end{pmatrix}. \tag{28}$$

We must diagonalize this matrix to obtain the right-handed neutrino eigenvalues and the Yukawa matrix in the basis of diagonal M_R. In analogy with the charged lepton and quark Yukawas we can expect $y_3 \gg y_2 \gg y_1$. Then we obtain, in the normal hierarchy case

$$M_{R_3} \simeq v^2 y_3^3 \frac{m_1^{-1}}{6}, \qquad M_{R_2} \simeq v^2 y_2^3 2m_3^{-1}, \quad \text{and} \quad M_{R_1} \simeq v^2 y_1^3 3m_2^{-1}. \tag{29}$$

Then, we have, $\tilde{m}_3 \simeq 6m_1$, $\tilde{m}_2 \simeq m_3$, and $\tilde{m}_1 \simeq \frac{2}{3} m_2$. This means that once again the right-handed neutrino that could dominate the energy density of the Universe is the heaviest one and its mass would be given by Eq. (25) with $\sin\theta_{13} \sim 1/\sqrt{6}$, i.e. a mass close or even above the Planck scale, unless y_3 is much smaller than 1.

Thus we see that with hierarchical Yukawa eigenvalues, similar to the up-quark eigenvalues, it is possible to have right-handed neutrino dominance of the energy density with consistent phenomenology, although only in rather fine-tuned situations where V_L is very close to the PMNS mixing matrix but not exactly equal. If we move to an extension of the minimal standard model with right-handed neutrinos (or minimal supersymmetric standard model) the same results can be obtained without such a tight fine-tuning.

From Eq. (16), we can see that besides the Yukawa mixing V_L we can still use different Yukawa eigenvalues. In a GUT with an underlying Pati-Salam symmetry, we would expect the neutrino Yukawa couplings to be related to the up-quark Yukawas, and in fact we would expect one of the neutrino eigenvalues of the order of the top Yukawa coupling [25]. However the two light Yukawa eigenvalues are less restricted. The masses of the right-handed neutrinos will depend on the Yukawa eigenvalues and we can make them as small as we wish. However, if y_3 is large, we will normally be in the same situation as before and the only neutrino with a sufficiently large lifetime will still be ν_{R_3}, which will be far too heavy. An interesting limit is when one of the Yukawa eigenvalues is exactly zero, $y_1 = 0$, and therefore, one of the light left-handed neutrino masses is zero. Again the simplest situation is when the right-handed Majorana matrix and the Yukawas, $Y_\nu^\dagger Y_\nu$ are simultaneously diagonalizable. In this case, it is clear that only two right-handed neutrinos will play a role in the seesaw mechanism and the third one will be completely decoupled from the seesaw. In fact this third neutrino does not couple to the doublets through Higgs Yukawa couplings. Given that right-handed neutrinos are singlets under the SM group, apparently these neutrinos do not decay at all. However, if we have a GUT symmetry, as for instance $SO(10)$ or a group containing $SU(2)_R$, at a high scale, the right-handed neutrino will decay with a lifetime

$$\begin{aligned}\Gamma_{\nu_R} &\simeq \alpha_{\rm GUT}^2 \frac{M_{R_i}^5}{M_{\rm GUT}^4} \\ &\simeq 1\times 10^{-18}(25\alpha_{\rm GUT})^2\left(\frac{M_{R_i}}{10^{10}\ {\rm GeV}}\right)^5 \\ &\quad\times\left(\frac{2\times 10^{16}\ {\rm GeV}}{M_{\rm GUT}}\right)^4\ {\rm GeV}. \end{aligned} \tag{30}$$

So, indeed we can see that this right-handed neutrino can dominate the energy density of the Universe if it is produced through another interaction, as a gauged B-L. In this case the production of the baryon asymmetry is not possible through gauge interactions. However, it is possible that this right-handed neutrino has nonvanishing complex Yukawa couplings when it unifies with the quarks at the GUT scale.[3] In this case the neutrino decay through GUT Higgses could violate CP and generate the observed baryon asymmetry.

Finally we would like to point out another possibility where we introduce more SM singlets mixed with the three "standard" right-handed neutrinos. An example of this situation is provided by the so-called "double seesaw" [27] mechanism. In this case, the right-handed Majorana masses are generated through a second seesaw with these additional singlets. The singlets would decay only through its mixings with the three right-handed neutrinos and, as we are introducing another free parameter, we can easily make any of these new singlets dominate the energy density of the Universe. Similarly, these singlets would easily generate the observed baryon asymmetry through the usual neutrino Yukawa couplings. The problem of how to generate a thermal abundance of these singlets could be solved again if they are charged under a gauged B-L symmetry.

IV. CONCLUSIONS

In this work, we have shown that an early epoch of matter domination by a long-lived massive particle can help to solve some of the problems of primordial inflation. We have seen that the large entropy production generated by the decays of such particle can dilute unwanted relics from higher temperatures, relaxing the constraints on the inflationary reheating temperature. In supersymmetric theories this mechanism can help to solve the gravitino problem. Moreover, a long period of matter domination reduces the number of e-foldings before the end of inflation at which the observable cosmological perturbations were generated. In the standard morel a natural candidate for such heavy, long-lived particle is a heavy right-handed neutrino. For a low enough mass of the lightest left-handed neutrino and neutrino Yukawa mixings sufficiently close to the PNMS mixing matrix, the right-handed neutrino dominates the energy density of the Universe for a long time and generates a large amount of entropy in its decay. In this case, we show that its decays can also generate the observed baryon asymmetry for right-handed neutrino masses well above the bound from gravitino overproduction.

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[3]For instance, we could think of a Georgi-Jarlskog vacuum expectation value distinguishing up-quark and neutrino Yukawas and being zero for the neutrinos [26].

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