Heavy Particles from Inflation

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We describe a simple and efficient mechanism by which very heavy particles are copiously created from a primordial inflationary epoch. It works for scalar fields which are massless or very light during inflation and acquire a large mass right after the end of inflation. Such particles can exist in realistic scenarios, as we illustrate with several examples in the context of both field and string theory. Long-wavelength fluctuations of these fields are generated during inflation with an almost scale-invariant spectrum and may give the dominant contribution to the energy density of the heavy fields at late times. Applications of our results to superheavy dark matter and leptogenesis are discussed.
1. Introduction

Inflation has become the standard paradigm for explaining the homogeneity and the isotropy of our observed Universe \([1]\) (for a recent review, see ref. \([2]\)). If at some primordial epoch the energy density of the Universe is dominated by some scalar field – the inflaton – whose kinetic energy is negligible, the corresponding vacuum energy gives rise to an exponential growth of the scale factor. During this phase a small, smooth region of size of the order of the Hubble radius grew so large that it easily encompasses the comoving volume of the entire presently observed Universe and one can understand why the observed Universe is homogeneous and isotropic to such high accuracy. Furthermore, it is now clear that structure in the Universe comes primarily from an almost scale-invariant superhorizon curvature perturbation. This perturbation originates presumably from the vacuum fluctuation, during the almost-exponential inflation, of some field with mass much less than the Hubble parameter \(H\) during inflation. Indeed, any scalar field whose mass is lighter than \(H\) suffers, in a (quasi) de Sitter epoch, quantum fluctuations whose power spectrum is independent of the length scales, if they are superhorizon \([3,4,5]\). On the contrary, quantum fluctuations of scalar fields whose mass \(M\) is larger than the Hubble rate during inflation are not efficiently excited, and their power spectrum is suppressed as \(\sim e^{-2M^2/H^2}\) \([6]\).

While massive fields may be produced after inflation, \textit{e.g.} during the process of preheating \([7]\), in this paper we show that there are interesting cases where fluctuations generated during inflation rather than quantum fluctuations produced after inflation give the dominant contribution to heavy particle production. The mechanism of heavy particle production from inflation we have in mind is based on a simple observation: particles which are massive in the present-day vacuum could have been very light during inflation. This implies that fluctuations of a generic scalar field, with mass \(M \ll H\) during inflation and \(M \gg H\) right after inflation, are copiously generated during inflation with an almost scale-invariant spectrum and become heavy and non-relativistic right at the end of inflation.
This provides an efficient mechanism to create massive particles from inflation without any exponential suppression. We will show that the number density of heavy particles at late times comes from the long-wavelength modes which are far outside the horizon at the end of inflation. As long as they are outside the horizon, these modes do not have a truly particle nature, but nevertheless may provide the dominant contribution to the energy density stored in the heavy field.

In sect. §2 we study the quantum creation of heavy particles from inflation. Next, we describe some interesting cases in which scalar fields can be massive in the present vacuum, but indeed massless during inflation: field-theoretical models are discussed in sect. §3 and D-brane models in sect. §4. We devote sect. §5 to possible applications of our findings. Finally, sect. §6 contains our conclusions.

2. Heavy particles from inflation

We start by considering a field $\chi$ which is nearly massless ($M < H$) during inflation, but very massive ($M \gg H$) right at the end of inflation. Our goal is to compute its abundance at the end of the inflationary stage and to show that the production of such massive field is very efficient.

Since the field $\chi$ is nearly massless during the inflationary epoch, its vacuum fluctuations in momentum space $\delta \chi_k$ are generated during the almost-exponential inflation. Indeed, the perturbations $\delta \chi_k$ acquire a nearly scale-invariant spectrum $[2,3,4,5]$. The fluctuations of the $\chi$ field have exponentially large wavelengths and – as long as their wavelength is larger than the horizon – for practical purposes they behave as a homogeneous classical field.

During inflation, the fluctuations of the $\chi$ field obey the equation

$$\ddot{\delta \chi}_k + 3 H \dot{\delta \chi}_k + \left(\frac{k}{a}\right)^2 \delta \chi_k = 0,$$  \hspace{1cm} (2.1)
where \(a\) is the scale factor and \(H\) is the Hubble rate. They can be conveniently described in terms of the variance

\[
\langle \chi^2 \rangle = \int \frac{d^3k}{(2\pi)^3} |\delta\chi_k|^2,
\]

which, during inflation, obeys the equation

\[
\frac{d\langle \chi^2 \rangle}{dt} = \frac{H^3}{4\pi^2}.
\]

In slow-roll models of inflation, the Hubble rate is not exactly constant, but slowly decreases with time according to the relation

\[
\dot{H} = -\epsilon H^2,
\]

where \(\epsilon = [V'(\phi)/V(\phi)]^2/(16\pi G)\) is one of the slow-roll parameters. Combining equations (2.3) and (2.4) and assuming that \(\epsilon\) does not change appreciably with time (in slow-roll, time derivatives of slow-roll parameters are proportional to the squares of slow-roll parameters and can be legally neglected), one finds

\[
\frac{d\langle \chi^2 \rangle}{dt} = -\frac{1}{8\epsilon \pi^2} \frac{dH^2}{dt}.
\]

If we indicate by \(H_*\) the initial value of the Hubble rate at the beginning of inflation, the value of the fluctuations of the massless field \(\chi\) depends on such an initial value as

\[
\langle \chi^2 \rangle = \frac{H_*^2}{8\epsilon \pi^2}.
\]

This means that the largest contribution to the variance of the \(\chi\) field is given by the time when the Hubble constant took its largest value or, equivalently, that the number density of \(\chi\) particles at late times comes predominantly from those infra-red long-wavelength modes which are far outside the horizon after inflation and crossed the horizon much before the end of inflation. As long as they remain outside the horizon these modes do not manifest a truly particle-like behaviour, but are indistinguishable from a classical homogenous field.
whose amplitude at the end of inflation is \( \sim \langle \chi^2 \rangle^{1/2} \sim 10^{-1} (H_*/\epsilon^{1/2}) \). This amplitude depends upon the model of inflation, but one can perform a simple estimate of it by assuming that the Universe expanded for at least \( \sim 60 \) e-folds before the end of inflation in order to solve the horizon and the flatness problems. In many models of inflation (but not in all) the slow-roll parameter \( \epsilon \sim 1/N \), when there are \( N \) e-folds to go till the end of inflation. Taking \( N \sim 60 \), one finds that the amplitude of the variance of the field \( \chi \) at the end of inflation is given by
\[
\langle \chi^2 \rangle^{1/2} \sim 0.9 H_* ,
\]
though inflationary models typically provide a number of e-folds much larger than 60 and the fluctuations of the \( \chi \) field at the end of inflation are expected to be greater.

In models of inflation where the inflaton field is well anchored at its false vacuum, such as the model of stringy old inflation proposed in ref. \([8]\), the Hubble rate is constant and the variance of the \( \chi \)-fluctuations is given by \( \langle \chi^2 \rangle = \frac{H^2}{4\pi^2} N \) thus providing a quantitative estimate similar to eq. (2.7).

Suppose now that the field \( \chi \) acquires a large mass \( M \) right at the end of inflation. This large mass is supposed to be greater than the Hubble rate at the end of inflation and at the beginning of the reheating stage. The long wavelength modes satisfy the equation
\[
\delta \dot{\chi}_k + 3H \delta \chi_k + M^2 \delta \chi_k \simeq 0 ,
\]
which tells us that modes beyond the horizon scale like \( \delta \chi_k \propto a^{-3/2} \cos Mt \). This means that the long wavelength modes of \( \chi \) become nonrelativistic and act like a classical homogeneous field. Their associated number density will be simply given by
\[
n_\chi = \frac{\rho_\chi}{M} = \frac{M\langle \chi^2 \rangle}{2} .
\]
The ratio at the end of inflation between the energy density stored in the heavy \( \chi \) particles and the energy density stored in the inflaton field is
\[
\frac{\rho_\chi}{\rho_\phi} \simeq \frac{1}{6\pi \epsilon} \left( \frac{M}{M_p} \right)^2 ,
\]
where $M_p = 1.2 \times 10^{19} \text{ GeV}$ is the Planck mass.

The energy density in the heavy nonrelativistic particles scales as $a^{-3}$, where $a$ is the scale factor. Meanwhile, during reheating the energy density stored in the inflaton field $\phi$ is redshifted away as $a^{-3}$ as well. This means that the initial ratio, given by eq. (2.10), does not change with time until the time when reheating occurs. At the time of reheating, the energy density stored in the inflaton oscillations gets converted into radiation with entropy density

$$s = 2\pi^2 g_\ast T_r^3/45,$$

where $g_\ast$ is the number of light degrees of freedom and $T_r$ is the reheating temperature. This yields

$$\frac{n_\chi}{s} = \frac{3T_r}{4M} \frac{\rho_\chi}{\rho_\phi} \approx \frac{MT_r}{8\pi\epsilon M_p^2}. \quad (2.11)$$

This ratio can be sizeable for large values of masses $M$ and small values of the slow-roll parameter $\epsilon$. The production of heavy particles which are massless during inflation and acquire a mass right after inflation can therefore be rather efficient.

Equation (2.11) is approximately valid also in the case of a particle with the same mass $M (\lesssim H)$ before and after inflation (see also ref. [9]). In this case, using the condition that $\chi$ is not thermally produced ($T_r \lesssim M$), we obtain an upper bound on its abundance,

$$n_\chi/s \lesssim H^2/(8\pi\epsilon M_p^2).$$

The present constraint on tensor perturbations ($H < 2 \times 10^{14} \text{ GeV}$ [10]), therefore requires $n_\chi/s \lesssim 1 \times 10^{-11}/\epsilon$. In the case under consideration in this paper, this limit can be violated, making use of the property that the mass $M$ after inflation can actually be much larger than $H$.

It is also interesting to compare the result of eq. (2.11), obtained for minimally-coupled scalars massless during inflation, with the case of conformally-coupled scalar particles or of fermions $\psi$, which have the same mass before and after inflation. Now the mass term is the only source of conformal-symmetry breaking, and such particles $\psi$ can be generated during the non-adiabatic transition between the end of inflation and the beginning of the radiation phase. Numerical simulations in chaotic inflation show that $\rho_\psi/\rho_\phi \simeq 10^{-3} H M_\psi/M_p^2$
when \( M_\psi \lesssim H \), while \( \rho_\psi / \rho_\phi \) is exponentially suppressed for \( M_\psi \gg H \). This gives

\[
\frac{n_\psi}{s} \simeq 10^{-3} \frac{HT_r}{M_p^2}.
\] (2.12)

Again, the requirement that \( \psi \) is not thermally produced \( (T_r \lesssim M) \) gives a stringent bound on its relic density \( n_\psi/s \lesssim 3 \times 10^{-13} (H/2 \times 10^{14} \text{ GeV})^2 \).

### 3. Massless states during inflation which become massive after inflation

In this section we want to show that it is indeed conceivable that states which are massive in the present-day vacuum could have been massless, or at least very light, during inflation. Let us suppose that inflation is driven by a scalar field \( \phi \) with potential \( V(\phi) \).

There are various possibilities for which a given field \( \chi \) is massless during inflation and very massive right at the end of inflation. We describe some of them, but the list is certainly not exhaustive.

The first option one can envisage is that the mass \( M \) of the \( \chi \) field is provided by an interaction term with the inflaton field itself of the form \( g^2 \phi^2 \chi^2 \), where \( g \) is some coupling constant. If during inflation the vacuum expectation value of the inflaton field is close to the origin in field space, \( \phi \simeq 0 \), the \( \chi \) field is effectively massless. However, as soon as inflation ends and the inflaton field rolls down the potential and acquires a large vacuum expectation \( v \), the field \( \chi \) becomes massive with mass squared \( M^2 = g^2 v^2 \gg H^2 \). This may happen, for instance, in the so-called small-field models [3,4] where the potential of the inflaton field, expanded around the origin, is given by

\[
V(\phi) = V_0 - \frac{m^2}{2} \phi^2 + \cdots .
\] (3.1)

Here \( V_0 \) represents the vacuum energy density driving inflation and the inflaton field slowly rolls down during inflation with a mass squared \( m^2 \ll H^2 \sim V_0/M_p^2 \). During inflation, the value of the inflaton field is exponentially close to the origin [3], \( \phi \sim \phi_f e^{-N(V_0/m^2M_p^2)} \),
where \( N \) is the number of \( e \)-folds to go till the end of inflation and \( \phi_f \) is the value of the inflaton field at the end of inflation. This means that the field \( \chi \) is practically massless during inflation and its vacuum fluctuations can be generated during the almost-exponential expansion epoch, as described in sect. \S \( 2 \).

Another possibility may arise in hybrid models of inflation \([12]\). The potential relevant for the inflationary trajectory is given by

\[
V(\phi, S) = \lambda (|S|^2 - \Lambda^2)^2 + \frac{m^2}{2} \phi^2 + h^2 |S|^2 \phi^2. \tag{3.2}
\]

The true vacuum of such a potential is given at \(|S| = \Lambda\) and \(\phi = 0\). However, if \(h^2 \phi^2 \gg 2\lambda \Lambda^2\), the curvature of the potential along the \(S\)-direction becomes positive and the potential reduces to \(V = \lambda \Lambda^4 + \frac{m^2}{2} \phi^2\). Inflation is driven by the vacuum energy density \(V_0 = \lambda \Lambda^4\) and ends when the inflaton field \(\phi\) reaches the critical point \(\phi_c^2 = 2(\lambda/h^2)\Lambda^2\). If the particle \(\chi\) acquires mass only through couplings with \(S\), it remains massless during inflation, when the vacuum expectation value of \(S\) vanishes, but it becomes massive after the field \(S\) has relaxed to its true minimum.

The mechanisms described above rely on tree-level potentials, but large masses for the scalar particles can be induced by radiative corrections. This is just another incarnation of the usual hierarchy problem of the Higgs potential, whose solution is still unknown. Supersymmetry or non-linearly realized symmetries could protect the scalar mass during inflation. A more exotic possibility is that \(\chi\) is the extra component of gauge fields in theories with more than four dimensions. If the corresponding gauge symmetry is unbroken during inflation, a mass term for \(\chi\) (which is a scalar field under the 4-dimensional Lorentz group) is forbidden. Once gauge symmetry is broken, after inflation, the \(\chi\) field becomes massive.

A second source of concern comes from induced interaction terms of the form

\[-(\xi/2)\sqrt{-g}R(g)\chi^2,\]

where \(R(g)\) is the scalar curvature, function of the metrics \(g\), and \(\xi\).
is a coupling constant. During inflation, this term generates a squared mass for $\chi$ equal to $6\xi H^2$, spoiling the condition $M \ll H$. This problem is actually generic in supergravity, since supersymmetric scalars are conformally coupled ($\xi = 1/6$). The $H^2$ terms in the $\chi$ mass can be made small, but at the price of a certain conspiracy between the different supersymmetry-breaking contributions. On the other hand, Goldstone bosons are minimally coupled ($\xi = 0$), and the non-linearly realized symmetry can simultaneously protect the $\chi$ mass from both quantum and $H^2$ corrections. This can be realized in our scenario if the inflaton background field breaks an exact (or approximate) global symmetry: a Goldstone (or pseudo-Goldstone) boson appears during inflation. If the inflaton field vanishes after inflation, the symmetry is restored and $\chi$ acquires a large mass.

Let us close this section with a final comment. In sect. §2 we have assumed that the classical field $\chi$ vanishes during inflation. There is no strong justification for this assumption in theories where the mass of the field $\chi$ is smaller than $H$ during inflation. However, the main purpose of the paper is to demonstrate that a large contribution to the energy stored in the heavy field $\chi$ may come from the quantum fluctuations created during inflation. In this respect, the computation of the energy density stored in the field $\chi$ presented in sect. §2 should be regarded as conservative.

4. Massless particles in $D$-brane models

Other set-ups where massless fields during inflation become massive after the inflationary epochs may arise in models of inflation inspired by supergravity and (super)string theories. In scenarios with $D$-brane inflation, light fields may arise as pseudo-Goldstone bosons associated to broken translational invariance [13]. In these models, the distance between a $D$-brane and an anti $D$-brane plays the role of the inflaton. If the branes are embedded in a string compactification where the background geometry has exact or approximate isometries, brane fluctuations transverse to the mutual distance may give rise to
almost massless fields. Inflation ends when the branes come close together and annihilate. In this process the light fields become very massive and decay. Indeed, after the end of the inflation the branes have disappeared and certainly cannot fluctuate anymore.

As a specific example, let us discuss the KKLMMT model [14] for D-brane inflation. The KKLMMT model is based on a compactification of the type IIB string theory on a Calabi-Yau manifold in the presence of orientifolds and fluxes for the antisymmetric tensors. As advocated in refs. [16,17], the fluxes can stabilize all the complex structure moduli of the compactification (the shape of the Calabi-Yau). The Kähler moduli (the volume) can be then stabilized by non-perturbative superpotentials generated by additional D7 branes (or euclidean instantonic D3-branes) naturally present in such compactifications [17,18]. The fluxes also induce a warp factor in the metric. The KKLMMT scenario

1 The KKLMMT model was proposed in order to overcome some of the problems of the D-brane inflation scenario in flat space but it is still plagued by the familiar $\eta$ problem of F-term inflation in supergravity (the brane position appears in the Kähler potential and the stabilization of the volume produces a mass for the inflaton of order $H$). Various ways to solve this problem, with or without fine tuning, have been proposed in ref. [14] and in subsequent works on the subject [15]. Here we take the conservative attitude of assuming that this problem has been solved and look for the presence of massless fields during the inflationary period.
corresponds to a Calabi-Yau geometry that develops a long warped AdS-like throat. The model may be thought as a stringy realization of the Randall-Sundrum model (RSI) [19], where the infra-red brane has been effectively regularized by an infra-red geometry and the ultra-violet brane has been replaced by the Calabi-Yau geometry (see Fig. 1).

The entire physics of inflation takes place in the throat region and it is mostly insensitive to the details of the ultra-violet Calabi-Yau region. In the KKLMMT an anti $D3$-brane is sitting at the infra-red end of the throat and a $D3$-brane is moving toward it. In the absence of anti branes, the $D3$ is a supersymmetric BPS object and it feels no force [20]. The only source of supersymmetry breaking in the compactification is the inclusion of the anti $D3$-brane. The insertion of a distant anti $D3$-brane just produces a very flat potential [14] for the scalar field corresponding to the position of the $D3$-brane, which plays the role of the inflaton field. The potential for the canonically normalized inflaton $\phi$ is

$$V(\phi) = T_3 a_0^4 \left( 1 - \frac{C}{\phi^4} \right),$$  

(4.1)

where $T_3$ is the tension of a $D3$ brane, $a_0$ is the redshift factor at the bottom of the throat and $C$ is a numerical constant whose value can be found in ref. [14]. If we ignore the order $H$ mass that the inflaton will acquire due to the Kähler potential [14], the condition of slow roll and the constraints from the density perturbations can be satisfied with reasonable values of the parameters\(^2\).

For our purposes, we notice that the KS solution [20] has an $SU(2) \times SU(2)$ isometry group. During its motion toward the infra-red, the $D3$-brane can fluctuate in the internal angular directions. Since during the inflationary period, the two brane set-up preserves the $SU(2) \times SU(2)$ symmetry, the scalar fields associated with the angular positions are almost massless and lighter than the Hubble rate. Furthermore, being angles, they do

\(^2\) In ref. [14] the integer fluxes $M, K$ are taken $\mathcal{O}(10)$, $T_3/M_\nu^3 \sim 10^{-3}$ and the warping is mild $a_0 \sim 10^{-4}$. The Hubble scale is consequently of typical order $10^9$ GeV.
not get masses from the volume stabilization mechanism. In the Kähler potential of four-dimensional supergravity the volume modulus $\rho$ always appears in the combination $\rho + \bar{\rho} - k(\phi_i, \bar{\phi}_i)$ \cite{21}, where $\phi_i$ collectively denote the position of the branes in the six internal directions, and $k$ is the Kähler potential for the geometry. This coupling generates a mass for the fluctuations $\phi_i$. However, since the geometry has isometries, $k(\phi, \bar{\phi})$ does not depend on some of the angles. In this way, the form of potential for the angular fluctuations is not affected by the stabilization mechanism and leaves some angles much lighter than the Hubble rate during inflation. At the end of inflation, the $D3$-brane comes close to the anti $D3$-brane. At stringy distances, we cannot ignore the fields coming from open strings connecting the branes. In particular, in a generic brane-antibrane system, the open string ground state is a tachyon $T$ with mass of order $M_s$, the string scale. Since the annihilation takes place at the bottom of the throat, the tachyonic mass is redshifted to $m_T^2 \sim -a_0^2 M_s^2$. We can write a toy-model potential for the inflaton $\phi$, the tachyon $T$ and the angular positions $\Theta_i$ which captures the salient features of the tachyonic condensation when the brane and anti-brane are close to each other

$$V(\phi, \Theta_i, T) = T_3 a_0^4 \left(1 - \frac{C}{\phi^4}\right) + T^2 (\phi^2 - a_0^2 M_s^2 + \tilde{V}(\Theta_i)) + \cdots,$$  \hspace{1cm} (4.2)$$

where the ellipses stand for higher-order power terms in the field $T$. In particular, the form of the potential accounts for the $SU(2) \times SU(2)$ sources of breaking when the two branes are close. Notice that this scenario is similar to that discussed in sect. §3 in the context of hybrid inflation. As long as $\phi$ is much larger that the (redshifted) string scale, $T$ has a large positive mass, its vacuum expectation value $\langle T \rangle$ vanishes and all the terms in the potential involving $T$ can be ignored. The inflationary potential is then independent of the angles $\Theta_i$. At stringy distances, $T$ becomes tachyonic and triggers the brane-antibrane annihilation. When $T$ condenses, all the fluctuation fields on the branes acquire masses roughly of order of the redshifted string scale. Taking into account that $a_0 \sim 10^{-4}$ at the bottom of the throat, these masses are $O(a_0 M_s)$ and much larger than $H \sim a_0^2 M_s$. All these massive fields will eventually decay.
There are two effects that break the $SU(2) \times SU(2)$ invariance and may prevent the angular fluctuations of the brane from being exactly massless during inflation. There is a small source of breaking of the $SU(2) \times SU(2)$ symmetry due to the presence of the anti $D3$-brane in the infra-red, which sits at a specific point on the three-sphere. This explicit breaking is though suppressed when the the distance between the two branes during the inflationary period is sizeable. The second effect is caused by the moduli stabilization that requires fluxes in the ultra-violet region typically breaking $SU(2) \times SU(2)$; it is well known indeed that a Calabi-Yau manifold has no isometries at all. Even this second effect is suppressed by the distance. An estimate based on the arguments in appendix A.2 of ref. [22] shows that the typical mass induces by the $SU(2) \times SU(2)$ breaking scales as $a^8(\phi)^3$, where $a(\phi)$ is the warp factor evaluated at the position $\phi$ of the brane. On the other hand, using formulae in appendix C of ref. [14], we can easily evaluate $a(\phi) \sim a_0^{2/3}$ when there are about 60 $e$-folds till the end of inflation. Therefore, the mass of the angular variables $m_{\Theta}^2 \sim a_0^{16/3}$ is suppressed compared to $H \sim a_0^4$ by powers of the redshift factor $a_0 \sim 10^{-4}$.

We finish this section by commenting on the dual interpretation of these massless fields. As familiar from the AdS/CFT correspondence and the holographic interpretation of the RS model, the throat part of the compactification can be effectively replaced by a strongly interacting gauge theory coupled to the four-dimensional gravity. In this picture, the $D3$ brane position corresponds to a flat direction in the moduli space of vacua of the gauge theory. The global symmetry $SU(2) \times SU(2)$ is spontaneously broken along this flat direction. The angular positions of the brane are identified with the massless Goldstone

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3 Actually, the smallest breaking effect could be of order $a^5(\phi)$ with a generic Calabi-Yau. As in ref. [22] we can however choose a manifold with a discrete $Z_2$ symmetry that forbids this effect.

4 The KS model with K and M fluxes is dual to a $SU(KM + M) \times SU(KM) N = 1$ supersymmetric theory with pairs of bifundamental fields $A_i$ and $B_i$, $i = 1, 2$, each one rotated by a global $SU(2)$ symmetry, with an $SU(2) \times SU(2)$ invariant superpotential. The same compactification with a $D3$ brane inserted can be represented by a gauge group $SU(KM + M + 1) \times SU(KM + 1)$.
bosons of the spontaneously broken global symmetry.

5. Applications of heavy particle production from inflation

In sect. §2 we have estimated the number density of heavy particles generated during inflation in the case in which the large masses are acquired after inflation. The generation mechanism does not directly involve the particle mass and it can be very efficient. We now proceed to describe some applications of our previous findings.

5.1. Superheavy dark matter

The case for dark nonbaryonic matter in the universe is today stronger than ever. The observed large-scale structure suggests that dark matter (DM) accounts for about 30% of the critical mass density of the universe \( \rho_C = 1.88 \times 10^{-29} \text{ g cm}^{-3} \) [23]. Despite this compelling evidence, the nature of the DM is still unknown. Some fundamental physics beyond the Standard Model (SM) is certainly required to account for the cold and slowly moving particles, \( \chi \), composing the bulk of the nonbaryonic dark matter. The most familiar assumption is that dark matter is a thermal relic, i.e., it was initially in chemical equilibrium in the early universe and its present-day abundance is inversely proportional to the annihilation cross section. Since the largest possible annihilation cross section is roughly \( M^{-2} \), one expects a maximum mass for a thermal DM particle, which turns out to be a few hundred TeV [24]. While a thermal origin for DM particles is the most common assumption, it is not the only possibility. DM particles may be produced independently of any microphysics process. It has been pointed out that DM particles might have never experienced local chemical equilibrium during the evolution of the universe and that their broken to \( SU(KM + M) \times SU(KM) \) along a flat direction where \( A, B \) acquire a VEV breaking the global symmetry \( SU(2) \times SU(2) \).
mass may be in the range $10^{12}$ to $10^{19}$ GeV, much larger than the mass of thermal relics $^{25,26,27,28,29}$. These superheavy DM particles have been called WIMPZILLAs $^{28}$. Independently of the production mechanism, the fraction of the total energy density of the Universe in superheavy DM particles today is given by

$$\Omega_\chi = \Omega_R \left( \frac{T_\chi}{T_0} \right) \frac{\rho_\chi}{\rho_\phi} \bigg|_* . \quad (5.1)$$

Here, $\Omega_R h^2 = 4.3 \times 10^{-5}$ is the fraction of critical energy density in radiation today, $T_0$ is the present temperature of radiation, and the subscript indicates the epoch at which heavy particles are generated. The present abundance of nonthermal DM particles is, as expected, independent of the cross section $^{25,26}$, and one can easily verify that, if it amounts to $\Omega_\chi \sim 1$, nonequilibrium during the evolution of the universe is automatic.

Several scenarios for superheavy DM particles have been developed, which involve production during different stages of the evolution of the universe. They may be generated in the transition between an inflationary and a matter-dominated (or radiation-dominated) universe due to the “nonadiabatic” expansion of the background spacetime acting on the vacuum quantum fluctuations. This mechanism was studied in details in Refs. $^{25,30}$, in the case of chaotic inflation. The distinguishing feature of this mechanism is the capability of generating particles with mass of the order of the inflaton mass (usually much larger than the reheating temperature, but at most of the order of the Hubble rate during inflation) even when the particles only interact extremely weakly (or not at all) with other particles, and do not couple to the inflaton. Superheavy DM particles may be also created during bubble collisions if inflation is completed through a first-order phase transition $^{31}$; at the preheating stage after the end of inflation with masses up to $10^{15}$ GeV if they are bosons $^4$; and up to the Planck scale if they are fermions $^{32}$; or during the reheating stage after inflation $^{27}$ with masses which may be as large as $2 \times 10^3$ times the reheat temperature. In addition, particles as heavy as the inflaton can be efficiently produced perturbatively, both in inflaton decay $^{33}$ and during thermalization of energetic inflaton decay products $^{34}$.
We can use the production mechanism described in sect, §2 and imagine that superheavy DM particles are massless during inflation and get their large mass right after inflation. In such a case, the present-day abundance is readily computed using eqs. (2.10) and (5.1) and we find \( \Omega_\chi h^2 = 0.1 \) for

\[
M \simeq \left( \frac{\epsilon}{10^{-2}} \right)^{1/2} \left( \frac{\text{TeV}}{T_r} \right)^{1/2} 4 \times 10^{12} \text{GeV}. \tag{5.2}
\]

The request that the DM particles are not produced by thermal processes \( (M \gtrsim T_r) \), implies that

\[
T_r \lesssim \left( \frac{\epsilon}{10^{-2}} \right)^{1/3} 3 \times 10^9 \text{GeV}. \tag{5.3}
\]

Therefore the proposed mechanism can account for DM in a large range of \( T_r \) which, in particular, can be compatible with an acceptable regeneration of thermal gravitinos. The upper bound on \( T_r \) reflects the extreme efficiency of the mechanism described in this paper in creating heavy particles from inflation.

5.2. Baryogenesis

Another natural application for the proposed mechanism of heavy-particle production is baryogenesis. Since baryon-number violation (or lepton-number violation, if sphalerons are active) is a necessary ingredient of baryogenesis, the relevant dynamics generally takes place at very short distances and involves very heavy particles. This is often at odds with the request of sufficiently low \( T_r \), coming from cosmological considerations on unwanted relics like, for instance, gravitinos or moduli (for a recent analysis, see ref. [35]). A mechanism for producing particles with masses much larger than \( T_r \) could resolve this conflict.

The situation is well illustrated by leptogenesis [36], one of the most concrete mechanisms to successfully reproduce the observed cosmic baryon asymmetry. Therefore we
will focus on this case, although our discussion can be extended also to the case in which baryogenesis occurs through the decay of some GUT field, nearly massless during inflation.

Thermal leptogenesis implies a lower bound on $T_r$, which depends on the right-handed neutrino mass, but it is always larger than about $10^9$ GeV \[34\]. If we want to consider lower values of $T_r$, we should rely either on variations of leptogenesis (soft leptogenesis \[38\], resonant leptogenesis \[39\], leptogenesis with sneutrino dominance \[40,41\]) or on non-thermal production of the right-handed neutrinos $N$. One possibility is to generate $N$ at reheating by direct inflaton decay \[42\]. Alternatively, the right-handed neutrinos could be produced during the reheating stage when temperatures much larger than $T_r$ are actually achieved \[27\]. If a preheating stage occurs, when large inflaton oscillations decay non-perturbatively, than a non-thermal density of heavy fermions can be generated \[32\]. Here we want to study if the mechanism proposed in this paper could also create a sufficiently large population of right-handed neutrinos.

Since $N$ are fermions, we cannot take direct advantage of fluctuations generated at inflation. The estimate given in eq. (2.12) shows that their density is too small to account for the baryon asymmetry. However, there are still some interesting possibilities for leptogenesis.

First, we can consider the scalar field $\phi$ whose lepton-number violating vev gives mass to the right-handed neutrinos though the coupling $\phi N e^T N$. If $\phi$ is massless during inflation and acquires a mass $M_\phi$ right after, it can play the role of the field $\chi$ described in sect. §2. In particular, this could happen if $\langle \phi \rangle$ is the dominant source of lepton violation during inflation; then the associated pseudo-Goldstone boson could obtain large fluctuations. The large $\phi$ number density given by eq. (2.11), after decay, will be transferred to $N$. Assuming that $\phi$ dominantly decays into $N$, the baryon asymmetry is given by

$$
\frac{n_B}{s} = a \epsilon_{CP} \frac{M_\phi T_r}{4\pi\epsilon M_p^2},
$$

(5.4)
Here $a = 0.35$ (both in the Standard Model and in minimal supersymmetry) relates the $B - L$ to the $B$ asymmetry through sphaleron interactions. For hierarchical neutrino masses, the CP asymmetry in $N$ decays ($\epsilon_{CP}$) has the upper bound \[5\]

$$\epsilon_{CP} < \frac{3 M_N m_{\nu}^{\text{atm}}}{16 \pi v^2} = 3 \times 10^{-7} \frac{M_N}{10^{10} \text{ GeV}},$$

where $M_N$ is the mass of the lightest right-handed neutrino. Requiring that eq. (5.4) reproduces the observed value $n_B/s = 8.7 \times 10^{-11}$ and using eq. (5.5), we obtain the constraint

$$\frac{T_r}{M_N} > \left( \frac{10^{15} \text{ GeV}}{M_N} \right)^2 \left( \frac{10^{18} \text{ GeV}}{M_\phi} \right) \left( \frac{\epsilon}{10^{-2}} \right) \times 10^{-4}. \quad (5.6)$$

This shows that, for heavy $N$, it is possible to generate the correct baryon asymmetry even if $T_r$ is several orders of magnitude smaller than $M_N$. This region of parameters ($T_r \ll M_N$) is not accessible to the standard thermal leptogenesis. The proposed mechanism is most efficient when $M_N$ is large, close to the GUT scale. Indeed, absence of thermal processes ($M_N \sim T_r$) gives a lower bound on the right-handed neutrino mass

$$M_N \gtrsim \left( \frac{10^{18} \text{ GeV}}{M_\phi} \right)^{1/2} \left( \frac{\epsilon}{10^{-2}} \right)^{1/2} \times 10^{13} \text{ GeV}. \quad (5.7)$$

Supersymmetry offers another possibility. The right-handed sneutrino could be the source of leptogenesis and, if its mass term is suppressed during inflation, its density before decay will be given by eq. (2.11). The same argument that lead us to eqs. (5.6) and (5.7) now gives us the constraints

$$T_r > \frac{\epsilon}{10^{-2}} \left( \frac{10^{15} \text{ GeV}}{M_N} \right)^2 2 \times 10^{14} \text{ GeV}. \quad (5.8)$$

$$M_N \gtrsim \left( \frac{\epsilon}{10^{-2}} \right)^{1/3} 6 \times 10^{14} \text{ GeV}. \quad (5.9)$$

These bounds are more stringent than in the previous case, where a larger density of right-handed neutrinos is achieved by considering very large $\phi$ masses.

\[4\] This bound can be relaxed if the neutrino-mass hierarchy is reduced \[14\].
6. Conclusions and discussion

In this paper we have demonstrated how a primordial epoch of inflation might be associated to the generation of very heavy particles. The mechanism is rather simple once it is realized that heavy states might have been either exactly or nearly massless during inflation. Long-wavelength fluctuations of such fields are generated during inflation with an almost scale-invariant spectrum and might provide the largest contribution to their energy density at late times. This mechanism is present not only in effective quantum field theories, but also in brane-world scenarios inspired by string theory where the lightness of would-be heavy states during inflation is guaranteed by exact or approximate isometries of the background geometry. We have described applications of our findings to superheavy dark matter and leptogenesis.

There are other interesting issues which would deserve further and careful investigation. First of all, the mechanism discussed in this paper might provide a natural way to implement the curvaton mechanism to generate cosmological perturbations. It has been proposed that the field responsible for the observed cosmological perturbations is some ‘curvaton’ field different from the inflaton [45]. During inflation, the curvaton energy density is negligible and isocurvature perturbations with a flat spectrum are produced in the curvaton field. After the end of inflation, the curvaton field oscillates during some radiation-dominated era, causing its energy density to grow and thereby converting the initial isocurvature into curvature perturbation. To be operative the curvaton mechanism requires the mass of the curvaton to be lighter than the Hubble rate during inflation and decay must occur after inflation. As we have shown, these requirements are satisfied, for instance, by brane models where the distance between a $D$-brane and an anti $D$-brane plays the role of the inflaton. Brane fluctuations transverse to the mutual distance give rise to almost massless fields during inflation, which become heavy and decay once inflation is terminated. The curvaton acts as a pseudo-Goldstone boson whose dynamics has been
Our findings might also be relevant for the dynamics of reheating after inflation. One can envisage two extreme cases. Suppose that the decay rate of the inflaton field into ordinary matter is highly suppressed. If so, populating the Universe with known particles might occur through $\chi$-states which are efficiently generated during inflation and very massive after inflation. On the contrary, if the inflaton couplings to ordinary matter are of order unity, defrosting the Universe after inflation is very efficient and might lead to high reheating temperatures $T_r$, possibly in conflict with the gravitino bound [35]. This drawback could be avoided if there is a subsequent release of entropy caused by the decay of heavy $\chi$-states abundantly produced during inflation.
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