

Review

# Cosmological Probes of Supersymmetric Field Theory Models at Superhigh Energy Scales

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**Abstract:** The lack of positive results in searches for supersymmetric (SUSY) particles at the Large Hadron Collider (LHC) and in direct searches for Weakly Interacting Massive Particles (WIMPs) in the underground experiments may hint to a super-high energy scale of SUSY phenomena beyond the reach of direct experimental probes. At such scales the supergravity models based on Starobinsky inflation can provide the mechanisms for both inflation and superheavy dark matter. However, it makes the indirect methods the only way of testing the SUSY models, so that cosmological probes acquire the special role in this context. Such probes can rely on the nontrivial effects of SUSY physics in the early Universe, which are all model-dependent and thus can provide discrimination of the models and their parameters. The nonstandard cosmological features like Primordial Black Holes (PBHs) or antimatter domains in a baryon-asymmetric universe are discussed as possible probes for high energy scale SUSY physics.

**Keywords:** supersymmetric models; supergravity; inflation; superheavy dark matter; Primordial Black Holes; cosmic antinuclei

## 1. Introduction

Studies of supersymmetric (SUSY) models in field theory and gravity involve proper combination of cosmological, astrophysical and experimental physical probes reflecting methods of cosmoparticle physics [1–6] that explores fundamental relationship of cosmology and particle physics. SUSY models provide solutions not only to the hierarchy problem, grand unification and the origin of Higgs mechanism, but also to the mechanisms of inflation, baryosynthesis and nonbaryonic dark matter.

Constraints on parameters of SUSY models inevitably involve cosmological probes. Such probes are usually supplementary to the search for SUSY particles at accelerators, specifying SUSY candidates for dark matter particles. It was expected that in parallel with the discovery of SUSY particles at the LHC, SUSY dark matter candidates could be found in direct dark matter searches, reproducing the Weakly Interacting Massive Particle (WIMP) miracle in the most expected and natural way. However, SUSY particles with the mass in the hundreds GeV range were not found at the LHC so far, as well as there seem to be no positive evidence for WIMPs in the underground dark matter experiments.

It may be the hint to a super-high energy scale of SUSY physics, at which direct probes are not available, making cosmological impact to be of special interest.

The fundamentals of SUSY models can naturally contain mechanisms for inflation and the origin of baryon asymmetry, thus providing the physical basis for inflationary cosmology with baryosynthesis and dark matter/energy which the modern cosmological paradigm is based on. Such mechanisms do involve super-high energy physics phenomena in the very early Universe, appealing to an extension of the set of theoretical tools to probe these phenomena (see Sections 2.2 and 2.3 for more).

In this paper, by following the general framework for describing the cosmological effects [7,8] and their possible applications to SUSY models [9] (see also [10–14]), we pay special attention to the non-standard features of the cosmological scenario, related to primordial nonhomogeneity of cosmological energy density and baryon excess. These features involving Primordial Black Holes (PBHs) formation and creation of antimatter domains can provide cosmological probes for SUSY models.

We start with a brief review of Starobinsky inflationary model and its embedding into supergravity, serving as the good example of viable basis for inflationary cosmology (Section 2). Then we consider the mechanisms of PBH formation (Section 3) specifying their links to possible first- and second order phase transitions in the early Universe. We consider observational probes for PBHs in (Section 4). Pending on the PBH mass these probes either use an analysis of effects of PBH evaporation (in which, in particular, production of superheavy gravitino is possible), or effects of PBH as the possible form of dark matter. A nonhomogeneous baryosynthesis based on SUSY models can give rise to antimatter stars in a baryon-asymmetric Universe (Section 5). We come to the conclusion (Section 6) that a detailed analysis of cosmological consequences of SUSY models can lead to their probes in the observed structure and evolution of the Universe.

## 2. Cosmological Traces of Starobinsky Supergravity Models

In this Section we briefly review Starobinsky inflationary model of  $(R + R^2)$  gravity, summarize the theoretical advantages of SUSY and supergravity, and very briefly describe the supergravity extensions of inflationary models in our approach, with inflaton belonging to a massive vector multiplet, see e.g., Ref. [15,16] for a review and the references therein.

### 2.1. Starobinsky Inflation

The possible gravitational origin of inflation is best described by the Starobinsky model defined by the action [17]

$$S_{\text{Star.}} = \frac{M_{\text{Pl}}^2}{2} \int d^4x \sqrt{-g} \left( R + \frac{1}{6m^2} R^2 \right), \quad (1)$$

where the reduced Planck mass  $M_{\text{Pl}} = 1/\sqrt{8\pi G_{\text{N}}} \approx 2.4 \times 10^{18}$  GeV, and the inflaton mass  $m$  have been introduced. This  $(R + R^2)$  gravity model is the simplest no-ghost extension of Einstein-Hilbert action in the context of the modified  $f(R)$  gravity theories defined by

$$S_f = \frac{M_{\text{Pl}}^2}{2} \int d^4x \sqrt{-g} f(R) \quad (2)$$

in terms of a function  $f(R)$  of the scalar curvature  $R$ . A generic  $f(R)$  gravity action (2) is classically equivalent to

$$S[g_{\mu\nu}, \chi] = \frac{M_{\text{Pl}}^2}{2} \int d^4x \sqrt{-g} [f'(\chi)(R - \chi) + f(\chi)] \quad (3)$$

with the real scalar field  $\chi$ , provided that  $f' \neq 0$  and  $f'' \neq 0$ . Here the primes stand for the derivatives with respect to a given argument. The classical equivalence is easy to verify because the  $\chi$ -field equation is algebraic and implies  $\chi = R$ . Next, the factor  $f'$  in front of the  $R$  in (3) can be eliminated

by Weyl transformation of metric  $g_{\mu\nu}$ , provided that  $f' > 0$ . As a result, the action (3) appears to be equivalent to the action of the scalar field  $\chi$  minimally coupled to Einstein gravity, while the scalar potential reads

$$V = \left( \frac{M_{\text{Pl}}^2}{2} \right) \frac{\chi f'(\chi) - f(\chi)}{f'(\chi)^2}. \quad (4)$$

The kinetic term of  $\chi$  gets the canonical normalization after a field redefinition  $\chi(\varphi)$  with

$$f'(\chi) = \exp\left(\sqrt{\frac{2}{3}}\varphi/M_{\text{Pl}}\right), \quad \varphi = \frac{\sqrt{3}M_{\text{Pl}}}{\sqrt{2}} \ln f'(\chi), \quad (5)$$

in terms of the canonical inflaton field  $\varphi$ . Then the total action takes the standard (quintessence) form

$$S_{\text{quintessence}}[g_{\mu\nu}, \varphi] = \frac{M_{\text{Pl}}^2}{2} \int d^4x \sqrt{-g} R - \int d^4x \sqrt{-g} \left[ \frac{1}{2} g^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi + V(\varphi) \right]. \quad (6)$$

The stability conditions in  $f(R)$  gravity theory are given by

$$f'(R) > 0 \quad \text{and} \quad f''(R) > 0, \quad (7)$$

respectively. The first condition above gets violated for  $R \leq -3m^2$ , which implies that the Starobinsky model can only be applied for  $R > -3m^2$ .

It is straightforward to derive  $R$  and  $f$  in terms of a scalar potential  $V$  from the equations above. One finds (see e.g., [18])

$$R = \left[ \frac{\sqrt{6}}{M_{\text{Pl}}} \frac{dV}{d\varphi} + \frac{4V}{M_{\text{Pl}}^2} \right] \exp\left(\sqrt{\frac{2}{3}}\varphi/M_{\text{Pl}}\right), \quad (8)$$

$$f = \left[ \frac{\sqrt{6}}{M_{\text{Pl}}} \frac{dV}{d\varphi} + \frac{2V}{M_{\text{Pl}}^2} \right] \exp\left(2\sqrt{\frac{2}{3}}\varphi/M_{\text{Pl}}\right). \quad (9)$$

These two equations define the function  $f(R)$  in the parametric form, i.e., the *inverse* transformation to (4). The classical equivalence (or duality) between the  $f(R)$  gravity theories (2) and the scalar-tensor (or quintessence) theories of gravity (6) is established provided the conditions (7) are satisfied.

In the case of Starobinsky model (1), one finds the celebrated potential

$$V(\varphi) = \frac{3}{4} M_{\text{Pl}}^2 m^2 \left[ 1 - \exp\left(-\sqrt{\frac{2}{3}}\varphi/M_{\text{Pl}}\right) \right]^2 \quad (10)$$

that is bounded from below, has a Minkowski vacuum at  $\varphi = 0$  and a plateau of positive height. The height is related to the inflationary energy density or the Hubble value during inflation. The plateau leads to slow roll of the inflaton field during inflation.

The duration of inflation is measured in the slow roll approximation by the e-foldings number defined by

$$N_e \approx \frac{1}{M_{\text{Pl}}^2} \int_{\varphi_{\text{end}}}^{\varphi_*} \frac{V}{V'} d\varphi. \quad (11)$$

Here  $\varphi_*$  is the inflaton value at horizon crossing, and  $\varphi_{\text{end}}$  is the inflaton value at the end of inflation. The end of inflation takes place when one of the slow roll parameters,

$$\varepsilon_V(\varphi) = \frac{M_{\text{Pl}}^2}{2} \left( \frac{V'}{V} \right)^2, \quad \eta_V(\varphi) = M_{\text{Pl}}^2 \left( \frac{V''}{V} \right), \quad (12)$$

approaches one. The reasonable value of  $N_e$  should be between 50 and 60, though it cannot be fixed independently so far, either in theory or by observations.

The Starobinsky model (1) is known as the excellent model of cosmological inflation, in very good agreement with the Planck 2018 data [19], as regards its predictions versus Planck measurements of the scalar perturbations tilt:  $n_s \approx 1 + 2\eta_V - 6\varepsilon_V \approx 1 - 2/N_e$  versus  $0.9649 \pm 0.0042$  (68% CL), and the tensor-to-scalar ratio:  $r \approx 16\varepsilon_V \approx 0.0037$  versus  $r < 0.064$  (95% CL), with the best fit at  $N_e \approx 55$ .

The Starobinsky model (1) is entirely based on gravitational interactions, so that it is truly geometrical. Its only parameter given by the inflaton mass  $m$  is fixed by the observed CMB amplitude as

$$m \approx 3 \cdot 10^{13} \text{ GeV} \quad \text{or} \quad \frac{m}{M_{\text{Pl}}} \approx 1.3 \times 10^{-5}. \quad (13)$$

Starobinsky inflation with the potential (4) can be compared to the other well studied (single-field) inflationary models, such as the mass term inflation with

$$V(\varphi) = \frac{1}{2} m^2 \varphi^2, \quad (14)$$

and the (pure) natural inflation [20] with

$$V(\varphi) = \Lambda^4 \left( 1 - \cos \frac{\varphi}{f} \right) \quad (15)$$

in terms of the two scales  $\Lambda \sim M_{\text{Pl}}$  and  $f \sim M_{\text{GUT}}$ . The mass term is the simplest option, whereas the natural inflation is motivated by the pseudo-Nambu-Goldstone nature of the inflaton-axion  $\varphi$ . However, both alternative proposals are either ruled out or are in tension with the Planck data, respectively.

## 2.2. Supersymmetry and Supergravity

SUSY is the leading candidate for new physics beyond the Standard Model of elementary particles. It is, therefore, quite natural to combine SUSY with gravity to supergravity. In addition, supergravity is the field theory of local supersymmetry that implies the general coordinate invariance. We consider only the minimal ( $N = 1$ ) supergravity because it is chiral. Chirality is known to be necessary for particle phenomenology and CP violation.

Supergravity has many attractive features:

- it has SUSY unifying bosons and fermions,
- it automatically includes General Relativity (GR),
- it is the conservative extension of GR and field theory (without violating their basic principles), which restricts a number of independent coupling constants,
- SUSY Grand Unified Theories (GUTs) gives rise to the perfect unification of electro-weak and strong interactions,
- the spectrum of matter-coupled supergravities with spontaneously broken SUSY has the natural candidate for dark matter particle, given by the Lightest SUSY Particle (LSP),
- SUSY can be used to stabilize the fundamental scales, towards solving the hierarchy problem,
- SUSY results in cancellation of quadratic UV-divergences in quantum loops,
- some supergravity theories can be considered as the low-energy effective actions of superstring theories, i.e., in quantum gravity.

It is, however, worth mentioning that SUSY is just a theoretical proposal that is not supported by any physical observations yet. But the current absence of SUSY evidence is not the evidence of the SUSY absence, because the scale of SUSY breaking is unknown and can be much larger than the electro-weak scale (this is called High Scale SUSY).

### 2.3. Starobinsky Inflation in Supergravity

A viable description of cosmological inflation and reheating in supergravity is non-trivial and highly constrained. In general, supergravity realizations of inflation necessarily give rise to more particles and fields. The standard approach assigns inflaton to a massive chiral multiplet with a Kähler potential  $K$  and a superpotential  $W$  (see e.g., Ref. [21] for a review and the references therein). Since inflation also leads to spontaneous SUSY breaking, it implies the existence of another (chiral) goldstino superfield. This setup generically leads to the so-called  $\eta$ -problem, so that substantial “theoretical engineering” is needed to avoid this problem and suppress the other (non-inflaton) scalars by some choice of  $K$  and  $W$ . It is possible to identify the inflaton and goldstino chiral superfields, in order to reduce a number of the physical scalars by half [22–25].

A supergravity extension of  $f(R)$  gravity was first proposed in the form of a chiral (F-type) action dubbed in the literature as the  $F(R)$  supergravity and leading to new interesting models of inflation, reheating and dark energy [18,26–34]. The chiral nature of  $F(R)$  supergravity action implies the absence of perturbative (loop) corrections due to the SUSY non-renormalization theorems. However, the  $F(R)$  supergravity cannot exactly reproduce the  $(R + R^2)$  gravity. The more general supergravity models were proposed with both D-type and F-type terms depending upon the chiral (scalar) curvature superfield and its conjugate, but they were also suffering problems with the extra propagating scalars [35,36].

When assigning inflaton to a massive *vector* multiplet, there is no extra scalar because the scalar field component of a real massive  $N = 1$  vector multiplet is also real. Goldstino can also belong to the same vector multiplet and can have the inflaton as its superpartner. The  $\eta$ -problem does not arise because the inflaton potential of a vector multiplet in supergravity is of the D-type instead of the F-type. The inflaton potential is given by a real function squared and allows any values for the cosmological tilts  $n_s$  and  $r$  [37,38]. However, the minima of this potential necessarily have the vanishing cosmological constant and the vanishing Vacuum Expectation value (VEV) of the auxiliary field  $D$ , so that only Minkowski vacua are allowed and SUSY is always restored after inflation. The relevant part of the supergravity Lagrangian reads

$$e^{-1}\mathcal{L} \supset -\frac{J_{CC}}{2\kappa^2}\partial_\mu C\partial^\mu C - \frac{g^2}{2\kappa^4}J_C^2, \quad (16)$$

where  $J = J(C)$  is arbitrary real function of the real (inflaton) scalar  $C$  that is the leading field component of the massive vector superfield, the  $J_C$  and  $J_{CC}$  denote the first and second derivatives of  $J$  with respect to  $C$ , respectively,  $g$  is the gauge coupling constant, and  $\kappa = M_{\text{pl}}^{-1}$ . The Starobinsky potential (10) is recovered by the choice

$$J = -\frac{3}{2\kappa^2}\ln(-\kappa C e^{\kappa C}), \quad \kappa C = -e^{\sqrt{2/3}\kappa\phi}, \quad g^2 = \frac{3}{2}\kappa^2 m^2. \quad (17)$$

The problem of spontaneous SUSY breaking after inflation in supergravity can be solved by adding Polonyi chiral superfield with a linear superpotential [39], defining the so-called Polonyi-Starobinsky (PS) supergravity. And the vector multiplet can get its mass via super-Higgs effect from a  $U(1)$  supersymmetric gauge field theory [40]. Then it becomes possible to envisage a unification of Starobinsky inflaton with the SUSY GUTs, similarly to the well known “flipped”  $SU(5) \times U(1)$  model arising from the compactified heterotic strings or from the intersecting D-branes.

However, in this case, Polonyi scalar gets mixed with the inflaton and causes instability of the inflationary trajectory. This problem can be solved by adding a new type of the Fayet-Iliopoulos term that does not require gauging the R-symmetry [41,42].

As was demonstrated in [43], dark matter in the PS supergravity can be represented by gravitino as the LSP particle, consistently with the High-Scale SUSY particle physics [44]. It was also found in [43] that Polonyi and gravitino particles can be efficiently produced during inflation. Perturbative decays of inflaton after inflation also produce Polonyi particles that rapidly decay into gravitinos. Along these

lines, a coherent picture of inflation and dark matter emerges, where the abundance of gravitinos produced after inflation fits the CMB constraints. We call this as the Super Heavy Dark Matter (SHDM) scenario with the gravitino mass of about  $10^{13}$  GeV. This scenario is based on the linearly realized supersymmetry and also avoids the notorious so-called gravitino and Polonyi problems related to the Big Bang Nucleosynthesis (BBN) and the DM overproduction.

#### 2.4. Cosmological Traces of Inflation and Reheating with SUSY

If R-parity is conserved, the Lightest Supersymmetric Particle (LSP) should be stable. Creation of these particles in the early Universe results in their survival and presence in the modern Universe in the form of dark matter. Therefore, the simplest cosmological effect of SUSY models is a primordial LSP abundance.

The maximal temperature of the inflationary Universe is the reheating temperature,  $T_r$ , after inflation. Hence, superheavy particles with the mass (We use the natural units  $\hbar = c = k = 1$ , unless they are specified otherwise)  $m \gg T_r$  cannot be in thermal equilibrium, and the gas of these particles cannot be originated from their freezing out in the early Universe. One should consider the detailed mechanism of their production in order to calculate their primordial abundance. Moreover in the case for the gravitino, the rate of its semi-gravitational interaction

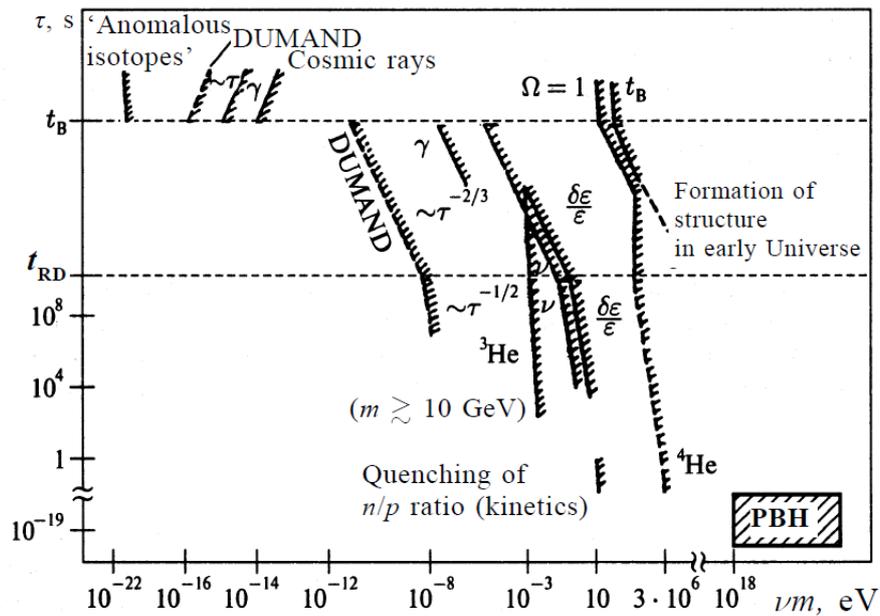
$$n\sigma v \sim T^3/M_{\text{Pl}}^2 \ll T^2/M_{\text{Pl}}$$

has never exceeded the rate of expansion at  $T < T_r \ll M_{\text{Pl}}$  so that gravitino could not be in equilibrium at any value of their mass.

Unstable particles, created in the early Universe, decay, if their lifetime is smaller than the age of the Universe,  $\tau < t_U$ . However, if their lifetime satisfies the condition  $\tau \gg (M_{\text{Pl}}/m) \cdot (1/m)$ , the presence of such particles in the Universe in the period  $t < \tau$  can lead to observable effects. These effects strongly depend on the modes of particle decay. Even completely invisible products of decay increase the contribution of the relativistic species in the cosmological density that is constrained by the data on light element abundance and the condition of the Large Scale Structure formation. The presence of photons or charged particles in the decay products leads to their contribution to the cosmic ray fluxes and gamma-ray background. If decay takes place earlier, when products of decay interact with matter and radiation such interaction influence the results of Big Bang Nucleosynthesis (BBN) or causes distortions of the CMB spectrum (see e.g., [9,45] and references therein). The set of these astrophysical probes for new particles with relative concentration  $\nu$  and mass  $m$ , decaying or present in the Universe at cosmological time  $t = \tau$ , is shown on Figure 1. It provides direct astrophysical test of existence of particles with lifetime  $\tau \geq 1$  s.

Test for existence of supermassive particles with lifetime shorter than 1 s is possible, if these particles dominate in the cosmological density before decay. Development of gravitational instability on the stage of their dominance is accompanied by formation of black holes. Spectrum of such primordial black holes is determined by particle mass, abundance and lifetime [46], while particle decay in the end of this stage produces additional entropy and thus influences the cosmological baryon asymmetry. In a multi-step analysis of cosmological consequences the prediction of superheavy, unstable but sufficiently long living new particles can be linked to observable astronomical effects.

Stability of new particles reflects the extension of the symmetry of the Standard Model. Such extensions involve new fundamental scales that determine the pattern of symmetry breaking and succession of cosmological phase transitions. Collision of bubbles, created in the first order phase transitions can produce black holes [47] (see [48,49] for a review). Symmetry of vacuum is changed in the second order phase transitions, giving rise to formation of topological defects, whose structure can be changed in successive phase transitions. It makes such structures unstable, but their existence can give rise to primordial non-homogeneous structures.



**Figure 1.** The constraints [6] on the possible effects of particles with relative concentration  $\nu$  and mass  $m$  at the cosmological time  $\tau$ .

### 3. Primordial Black Holes from Superhigh Energy Physics

Primordial Black Holes (PBHs), first suggested by Zeldovich and Novikov [50], provide a very sensitive probe of cosmology of the very early Universe.

Pending on their mass PBHs can survive to the present time, being a specific form of dark matter (see e.g., [51]), or evaporate by Hawking mechanism [52], leaving observable effects of products of their evaporation. If predicted effect of PBHs contradicts the observational data, it provides constraints on the physics of the very early Universe underlying the mechanisms of PBH formation. On the other hand, effects of PBHs and their evaporation add new cosmological parameters that can help to solve some existing astrophysical problems.

Formation of a black hole implies metric fluctuation of the order one and such a strong inhomogeneity is highly improbable. For the equation of state  $p = \gamma\epsilon$ , with  $0 \leq \gamma \leq 1$ , the probability of black hole formation from a fluctuation with the amplitude  $\delta$  is given by [53]

$$W_{PBH} \propto \exp\left(-\frac{\gamma^2 \delta^2}{2 \langle \delta^2 \rangle}\right), \tag{18}$$

where the dispersion is  $\langle \delta^2 \rangle \ll 1$ . Therefore, the probability of a fluctuation  $\delta \sim 1$  is exponentially small. It makes the PBH spectrum exponentially sensitive to the softening of the equation of state ( $\gamma \rightarrow 0$ ), corresponding to early dust-like stages, or to the increase of density fluctuations ( $\langle \delta^2 \rangle \rightarrow 1$ ), corresponding to much stronger nonhomogeneity of the early Universe, as compared with expected from the flat Harrison-Zeldovich spectrum. These effects can be cosmological tracers of superhigh energy physics.

We review in this Section following [7–9,51], the relationship of the corresponding mechanisms of PBH formation and parameters of new physics.

#### 3.1. PBHs from Superheavy Metastable Particles

At the radiation dominated stage after reheating, particles with the mass  $m$  and frozen-out or frozen-in concentration  $n$  must dominate over relativistic species with concentration  $n_r$  at  $T < T_0 = rm$ , where  $r = n/n_r$ . To avoid contradictions with the observational data such early dust-like stage should

end by particle decay long before the period of Big Bang Nucleosynthesis, so that the particles should have lifetime  $\tau \ll 1$  s and the stage of their dominance takes place in the period  $t_0 < t < \tau$ , where  $t_0 = M_{\text{Pl}}/T_0^2$ . It was first noticed in [54] that at the sufficiently long early dust-like stage formation of gravitationally bound systems should be inevitably accompanied by the formation of black holes that remain in the Universe after particle decay. Spectrum of these PBHs is determined by particle mass, primordial abundance and lifetime.

At the dust-like stage density fluctuations grow as

$$\delta(t) = \frac{\delta\rho}{\rho} \propto t^{2/3}.$$

Fluctuations that enter horizon at  $t = t_i > t_0$  with amplitude  $\delta(t_i)$  decouple from the cosmological expansion and form gravitationally bound systems with  $\delta(t_f) \sim 1$ , at

$$t_f = t_i \delta(t_i)^{-3/2}. \quad (19)$$

Such systems can form black-holes either directly in the course of contraction after an especially homogeneous and isotropic configuration separates from the expansion or in the result of evolution of gravitationally bound objects. In the latter case the rate of evolution is determined by particle properties. The timescale of evolution of gravitationally bound system of collisionless particles is at best of the order of  $N^{2/3}t_f$ , where  $N$  is the total number of particles in the system. However, if the particles are coupled to relativistic species (as it is the case for magnetic monopoles) energy loss by radiation accelerates evolution of the object making the timescale of the black hole formation of the order of  $t_f$  [55,56].

The minimal probability to form PBH at the dust-like stage is given by probability of direct PBH formation which is determined by a probability of realization of especially homogeneous and isotropic configurations. In the essence the idea of [54] was to find a probability that the configuration is so homogeneous and isotropic that it contracts within its gravitational radius, when it separates from the cosmological expansion. For  $\langle \delta^2 \rangle \ll 1$  such probability is suppressed  $W \propto \langle \delta^2 \rangle^{13/2}$  [8,54].

Direct mechanism gives rise to formation of spectrum PBH masses in the interval

$$M_0 \leq M \leq M_{bhmax}.$$

Here  $M_0$  is the mass within the cosmological horizon at time  $t \sim t_0$ , equal to [8,54]

$$M_0 = \frac{4\pi}{3} \rho t_0^3 \approx M_{\text{Pl}} \left( \frac{M_{\text{Pl}}}{rm} \right)^2. \quad (20)$$

The maximal mass  $M_{bhmax}$  is indirectly determined by the condition that fluctuations at this can manage to grow, separate from expansion and collapse before particles decay at  $t = \tau$ . If the spectrum is scale-invariant  $\delta(M) = \delta_0$ , the maximal mass is given by [8]

$$M_{bhmax} = M_{\text{Pl}} \frac{\tau}{t_{\text{Pl}}} \delta_0^{-3/2} = M_{\text{Pl}}^2 \tau \delta_0^{-3/2}. \quad (21)$$

The direct mechanism of PBH formation is independent of the nature of particles, dominating at the dust-like stage. Therefore it can also take place, if preheating stage of inflaton field oscillations is dust-like [57].

### 3.2. PBHs from Phase Transitions during Inflation

The probability of PBH formation can be enhanced, if the amplitude of small scale fluctuations is much higher, than it is extrapolated from the amplitude at the large scales, deduced from the CMB and LSS data. Though this amplitude tends to slightly decrease to smaller scales in the CMB and

LSS measurements, this trend cannot be proved at the galactic and subgalactic scales, at which the Universe is strongly nonhomogeneous. It makes principally possible to admit an enhancement of primordial nonhomogeneity at the scales, smaller than the scale of galaxies.

Physical motivation for such a possibility comes from the following argument. Realistic particle physics model doesn't reduce its prediction to the candidate of inflaton but should also contain prediction of other scalar fields that can influence the form of density fluctuation spectrum.

In particular, it was first pointed out in [58] that phase transitions during inflation can lead to spikes in the spectrum due to interaction of a Higgs field  $\phi$  with the inflaton  $\eta$ . Such interaction induces in the Higgs potential a positive mass term  $+\frac{v^2}{2}\eta^2\phi^2$  that changes sign in the course of the inflaton slow rolling

$$V(\phi, \eta) = -\frac{m_\phi^2}{2}\phi^2 + \frac{\lambda_\phi}{4}\phi^4 + \frac{v^2}{2}\eta^2\phi^2 \quad (22)$$

at the critical value of the inflaton amplitude  $\eta_c = m_\phi/v$ . It leads to spikes in the spectrum of density perturbations. The probability of PBH formation is strongly increased [9,59], when these spike-like perturbations reenter the horizon. The mass of PBH depends exponentially on the e-folding, corresponding to the phase transition, so that not only small mass, but even stellar and superstellar mass PBHs can be formed by this mechanism [8,9,51].

### 3.3. PBHs from Bubble Collisions in First-Order Phase Transitions

If transition from false to true vacuum proceeds by bubble nucleation as the first order phase transition, black holes can be formed in the collisions of expanding true vacuum bubbles (see [8,9] for review and reference). Mutual penetration of two bubble walls causes restoration of false vacuum in the region of collision and after casual separation of this region from the walls a false vacuum bag is formed [60]. Contraction of this bag within gravitational radius leads to formation of a black hole [48]. The mass of this PBH is given by (see [48])

$$M_{BH} = \gamma_1 M_{bub} \quad (23)$$

where  $\gamma_1 \simeq 10^{-2}$  and  $M_{bub}$  is the mass within the bubble volume released in the false vacuum decay.

If first-order phase transition takes place in the end of inflation, in the percolation regime true vacuum bubbles are of Hubble size. According to [48], collisions of such bubbles should lead to copious production of PBHs with masses given by Equation (23) with  $M_{bub}$  equal to the mass within the Hubble horizon at the end of inflation

$$M_{bub} \approx \frac{M_{Pl}^2}{H_{end}} \quad (24)$$

For the Hubble constant in the end of inflation  $H_{end} \approx 4 \times 10^{-6} M_{Pl}$  the mass of PBHs is of the order of  $M_0 \approx 1$  g. The contribution of these PBHs into the total density in the period of their formation was estimated in [48] as  $\beta_0 \approx 6 \times 10^{-3}$ . Such PBHs should evaporate at  $t \sim 10^{-27}$  s and the Hawking temperature of their evaporation is of the order of  $10^{13}$  GeV.

### 3.4. Massive Primordial Black Holes from the Succession of U(1) Phase Transitions

If global U(1) symmetry is broken spontaneously at large energy scale  $f$ , it results in a continuous degeneracy of vacua. Then successive explicit symmetry breaking at a smaller energy scale  $\Lambda \ll f$  leads to change of degeneracy of vacuum from continuous to a discrete one.

If the first phase transition proceeds at the inflationary stage, and the second one after reheating, formation and collapse of closed domain walls leads to formation of PBHs. Their mass is determined by the parameters  $f$  and  $\Lambda$  and can reach the values of stellar and superstellar mass up to the mass of AGNs [8,9,61]. Such black holes can form massive PBH clusters [49,62–64] and shed new light on the problem of galaxy formation [49,64,65].

Massive wall formation is illustrated by a model of complex scalar field  $\varphi = re^{i\theta}$  with the potential [9,49,61–63]

$$V(\varphi) = \lambda(|\varphi|^2 - f^2/2)^2 + \delta V(\theta). \quad (25)$$

This potential implies  $U(1)$  symmetry breaking at the stage of inflation and the form of the complex field  $\varphi \approx f/\sqrt{2} \cdot e^{i\theta}$ . Renormalization of the Lagrangian by instanton effects (see, e.g., [9]) generates an additional term

$$\delta V(\theta) = \Lambda^4 (1 - \cos \theta) \quad (26)$$

with  $\Lambda \ll f$ . This term is negligible at the inflationary and postinflationary stages. It becomes significant at the temperature  $T \sim \Lambda$ . Then the quantity  $f\theta$  acquires the meaning of a massive field with the mass

$$m_\theta = \frac{2f^2}{\Lambda}. \quad (27)$$

If the phase is  $\theta_0 < \pi$  at the stage of inflation, corresponding to the currently observed part of the Universe, at each successive step of inflation it fluctuates in smaller scales as

$$\delta\theta = H/2\pi f, \quad (28)$$

where  $H$  is the Hubble constant in the period of inflation and at some step such fluctuations can make  $\theta > \pi$  in some region, much smaller the currently observed part of the Universe. This domain of  $\theta > \pi$  is surrounded by the regions with  $\theta < \pi$ . At the successive steps after first crossing  $\pi$  the phase can fluctuate back to  $\theta < \pi$  within the region with  $\theta > \pi$  and in the nearby regions with  $\theta < \pi$  to new smaller domains with  $\theta > \pi$ . When the term (26) becomes significant and continuous degeneracy of vacuum is changed by discrete values of potential minimum at  $\theta_{vac} = 0, \pm 2\pi, \pm 4\pi, \dots$  at the border of regions with  $\theta < \pi$  and  $\theta > \pi$  appear closed domain walls, separating domains with vacuum values  $\theta_{vac} = 0$  and  $\theta_{vac} = 2\pi$  (see [8,9,66] and references therein).

Evolution of the structure of such domain walls with the width  $\sim 1/m \sim f/\Lambda^2$  and surface energy density  $\sim f\Lambda^2$ , studied in details in [61], can lead to formation of PBH clusters (see e.g., [9] for review and references).

The mass range of PBHs, formed in the collapse of closed walls, is determined by the parameters  $f$  and  $\Lambda$ . The principally maximal PBH mass follows from the condition that before it enters the cosmological horizon the wall does not dominate locally. It is given by [49]

$$M_{max} = \frac{M_{Pl}}{f} M_{Pl} \left(\frac{M_{Pl}}{\Lambda}\right)^2. \quad (29)$$

PBH with the locally maximal mass  $M < M_{max}$  is surrounded by a cloud of PBHs with smaller mass. Their minimal mass, which is determined by the condition that the gravitational radius of wall exceeds its width, is equal to [49,62]

$$M_{min} = f \left(\frac{M_{Pl}}{\Lambda}\right)^2. \quad (30)$$

Collapse of wall can be accompanied by gravitational wave (GW) signal. Its features need special studies. Qualitative analysis of [8] shows that the corresponding primordial GW spectrum is peaked at the frequency  $\nu_p$

$$\nu_p = 3 \cdot 10^{11} (\Lambda/f) \text{ Hz} \quad (31)$$

and the energy density of this primordial GW background can be up to

$$\Omega_{GW} \approx 10^{-4} (f/M_{Pl}). \quad (32)$$

It follows from Equation (32) that the contribution in the total density of this primordial GW background can reach  $\Omega_{GW} \approx 10^{-9}$  for  $f \sim 10^{14}$  GeV. At this value of  $f$  the peak frequency  $\nu_p$  depends on the value  $\Lambda$ , which can be in the range [8]

$$1 < \Lambda < 10^8 \text{ GeV}, \quad (33)$$

corresponding to the maximum of the spectrum at

$$3 \times 10^{-3} < \nu_p < 3 \times 10^5 \text{ Hz}. \quad (34)$$

Formation of PBHs in clusters facilitates formation of their binaries and their coalescence can be another source of gravitational wave signals [9,64]. These effects can be important for interpretation of the results of gravitational wave experiments.

#### 4. Observational Probes for PBHs

##### 4.1. Gravitino Production by PBH Evaporation

PBHs with the mass smaller than  $10^{14}$  g do not survive to the present time, but effect of their evaporation, revealed by S. Hawking [52], can lead to observational constraints similar to constraints on unstable particles, presented on Figure 1. Due to evaporation, PBH becomes a black-body source of particles with temperature

$$T_{eva} = \frac{1}{8\pi GM} \approx 10^{13} \text{ GeV} \frac{1 \text{ g}}{M}. \quad (35)$$

The time scale of evaporation has the form (see e.g., [8,9])

$$\tau_{BH} = \frac{M^3}{g_* M_{Pl}^4} \quad (36)$$

where  $g_*$  is the number of evaporated species with the account for their statistical weight.

However, the constraints on evaporating PBHs are not simply reduced to constraints on unstable particles. Indeed, independent of the strength of its interaction, any particle with mass  $m < T_{eva}$  can be produced by the gravitational field of the evaporating black hole. It provides a mechanism for copious production of superweakly interacting particles.

The observational constraints on the relative contribution  $\alpha(M)$  of density of PBH of mass  $M$ ,  $\rho_{PBH}$ , into the total density  $\rho_{tot}$  in the period of PBH evaporation

$$\alpha(M) \equiv \frac{\rho_{PBH}(M)}{\rho_{tot}} \quad (37)$$

provides the constraint on the contribution  $\beta(M)$  of PBHs in the period of their creation. At relativistic dominant (RD) stage of expansion the total density is  $\propto a^{-4} \propto T^4$ , while PBH contribution in the total density depends on  $a$  as  $a^{-3} \propto T^3$ . It converts even modest constraints on  $\alpha(M)$  into very strong constraints on  $\beta(M)$ . For instance, if PBHs are formed and evaporate at the RD stage, the constraint on  $\beta(M)$  is given by [46,67]

$$\beta(M) = \alpha(M) \frac{M_{Pl}}{M}. \quad (38)$$

At matter dominated stages the contribution of PBHs into the total density doesn't grow. The existence of such stages in the period between the times of formation and evaporation of PBHs should be taken into account in the constraint on  $\beta(M)$  and makes it correspondingly weaker (see [8,46,68,69] for details). Therefore analysis of allowed probability of PBH formation strongly depends on the details of a cosmological scenario, involving a possibility of such early dust-like stages.

Copious production of stable gravitino in PBH evaporation [70] makes possible to trace the existence of even low mass PBHs, evaporating in the very early Universe at  $t \ll 1$  s [71]. In particular, if inflationary stage ends by a first order phase transition, formation and successive evaporation of miniPBHs with mass 1 g should result in overproduction of stable gravitino with mass  $10^{13}$  GeV, excluding such scenario.

Indeed, such miniPBHs have the temperature of evaporation  $T_{eva} = 10^{13}$  GeV and the mass fraction of evaporated gravitino is of the order of  $1/g_*$ . Here the number of evaporated species  $g_* \sim 10^2$  takes into account their statistical weight. For the probability of PBH formation  $\beta_0 \geq 10^{-3}$  estimated in Section 3.3, in the period of PBH evaporation gravitino contribution  $\rho_G$  into the total density  $\rho_{tot}$  is of the order of

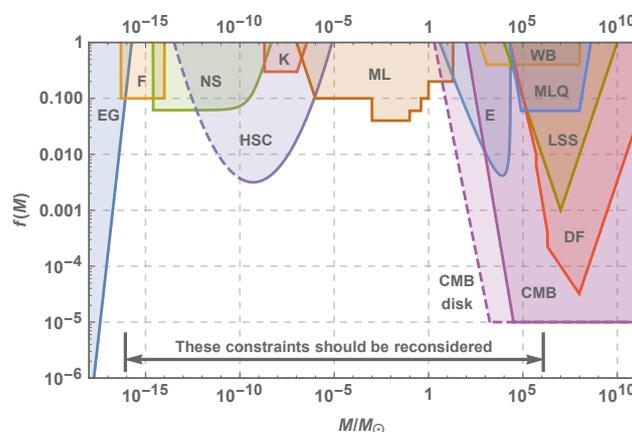
$$\frac{\rho_G}{\rho_{tot}} \sim \frac{\beta_0}{g_*} \geq 10^{-5}. \quad (39)$$

If there is no additional inflationary stage after PBH evaporation that can suppress gravitino abundance, stable superheavy gravitino should start to dominate at the RD stage in the contradiction with the observational data.

Adjustment of parameters of Starobinsky supergravity scenario can make PBH spectrum not only compatible with the observational data, but even provide mechanism for freeze in of supermassive gravitino as the dominant form of cosmological dark matter.

#### 4.2. PBH Dark Matter

PBHs with the mass exceeding  $10^{14}$  g should survive to the present time and contribute to the modern dark matter density. Critical analysis by [64] of constraints [72] on relative contribution of PBHs into dark matter density is presented on Figure 2. It is noted that PBH formation in clusters can influence the constraints [72] excluding PBH dominance in dark matter. On Figure 2 EG denotes constraints from observation of extragalactic gamma-ray background, while F are from femtolensing and NS from destruction of neutron stars. The constraints from searches for gravitational microlensing events by Subaru Hyper Suprime-Cam are denoted as HSC. The line K corresponds to constraints from the data of Kepler satellite. The line ML comes from MACHO. The exclusion line MLQ is from quasars millilensing. The constraints from wide binary are given by line WB. The line E corresponds to constraints from destruction of star cluster in Eridanus dwarf galaxy. Line WMAP is obtained from CMB distortion due to accretion effects. Constraints from effects of dynamical friction in our Galaxy are shown by line DF. The line LSS corresponds to constraints from large-scale structure. CMB constraint was considered in [64] both for case of spherically symmetric accretion [73] and for the stronger constraint (“CMB disk”) one [72] (shown by dashed line).



**Figure 2.** The critical analysis by [64] of constraints on Primordial Black Holes (PBH) dark matter contribution into the total density by [72] with the account for the less stringent CMB constraints by [73].

The mass range, in which the constraints should be re-considered with te account for PBH clusters, is shown by double-headed arrow. It indicates that PBH dominance in the dark matter density may not be excluded for a rather wide range of PBH masses.

Even subdominant component of massive PBH clusters can lead to important consequences for interpretation of gravitational wave signals from coalescence of massive black holes. The latest GWTC catalog [74] contains now 11 events of coalescence of compact binaries, with only one event of neutron star-neutron star binary coalescence. All the other 10 events correspond to coalescence of binary systems of black holes with mass, exceeding 10–20 Solar masses each. Formation of such systems in the evolution of massive stars is rather problematic, while it is rather natural for clusters of massive PBHs, in which formation of binaries is much more probable as compared with the case of random PBH distribution. Repeating merging of black holes in clusters may be another signature for massive PBH clusters [64,75]. The existing statistics is evidently not sufficient to make any definite conclusion on this possibility. However, repeating detection of four GW signals in the August of 2017 may be an interesting hint to such a possibility.

### 5. Anti-Matter Stars in Our Galaxy?

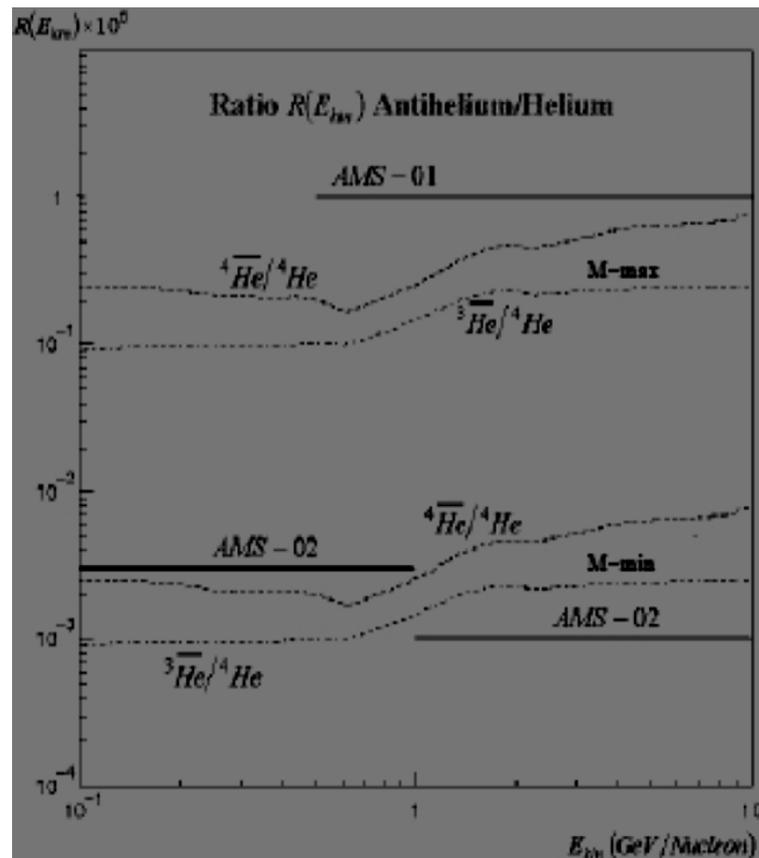
Any mechanism of generation of baryon excess in the baryon-symmetrical Universe can in the case of strong nonhomogeneity of baryosynthesis lead to generation of antibaryon excess [76]. The fate of such antimatter domains depends on their size. Their survival can lead to appearance of antimatter stars in baryon asymmetrical Universe.

Implementation of the model of spontaneous baryosynthesis [77] (see review in Reference [78]) has provided quantitative analysis of distribution of antimatter domains and originated from nonhomogeneous baryosynthesis in the inflationary Universe [79]. It was shown for a reasonable choice of parameters that the predicted number of antimatter domains of globular cluster scale (with masses of several thousand Solar masses) can be comparable with the number of observed galaxies, while probability of antimatter galaxies and antimatter domains at larger scales is strongly suppressed, so that there is no significant effect of annihilation at the borders of large domains in the observed gamma ray background [79].

The minimal survival scale of antimatter domains is about  $10^3 M_{\odot}$ , being of the order of the minimal mass of globular clusters. In our Galaxy such domain can form an antimatter globular cluster [80]. Indeed, globular clusters are old objects, and if antimatter cluster is formed before its protocloud meets and annihilates with matter clouds, successive annihilation of matter gas can take place only on the surface of antimatter stars. Since globular clusters are distributed in the galactic halo, where the matter gas density is low, such annihilation makes antimatter globular cluster a vary faint gamma ray source. Owing to a relatively small gravitational potential of globular cluster, antimatter gas escapes the cluster, polluting galactic halo by antimatter, in which anti-hydrogen (antiprotons) dominate. Antiproton annihilation in the galactic disc can reproduce [81] the observed gamma ray background. However, this background can have ordinary astrophysical origin.

Search for pieces of antimatter like antimatter meteorites [82] could be much stronger evidence favoring its macroscopic galactic origin. The existence of antihelium component in cosmic rays [83] is rather stable prediction of hypothesis of antimatter globular cluster [80]. Production of anti-He-3 and anti-He-4 in cosmic rays is strongly suppressed, while antimatter nucleosynthesis should make antihelium-4 the second antimatter element after antihydrogen by its abundance. Evolution of antimatter main sequence stars leads to additional production of antihelium, while propagation of heavy anti-nuclei in the matter gas is accompanied by their annihilation in which antihelium fragments are copiously produced. The estimated cosmic antihelium flux is shown on Figure 3.

This prediction is not only accessible to searches for cosmic ray antinuclei in the AMS02 experiment, but also provide decisive experimental test for antimatter globular cluster hypothesis.



**Figure 3.** The expected flux of antiHelium-3 and antiHelium-4 [49] from antimatter globular cluster in our Galaxy.

AMS collaboration [84,85] continuously presents results of these searches. There are about ten clear candidates for antihelium-3. Two events may be interpreted as antihelium-4. It is expected that more statistics will be available to 2024 and the result with the significance level of 5 standard deviation can be obtained. The expected progress will shed light on the reasons of dominance of events with anti-helium-3 over anti-helium-4, as well as of the absence of candidates for anti-deuterium in the presented results. In any case antihelium-3 events cannot be explained as the secondary effect from cosmic ray interactions [86]. Therefore their confirmation would be a strong argument, favoring antimatter globular cluster hypothesis.

Formation of dense antistars and their possible search were considered in [87] within an extension of the supersymmetric Affleck-Dine mechanism of baryosynthesis. It would be interesting to check, whether implementation of Starobinsky supergravity can provide nonhomogeneous baryosynthesis and reproduce hypothesis of antimatter globular cluster in our Galaxy.

## 6. Supersymmetry in the Context of Cosmoparticle Physics

The development of SUSY models towards superhigh energy scales makes their cosmological consequences to be of special interest. In the lack of an opportunity for searching supersymmetric particles at the LHC and SUSY dark matter candidates in the underground experiments, an experimental test of supersymmetric models becomes the matter of astroparticle physics and cosmology. Cosmological probes grow from supplementary tools to the main source of experimental information.

With such development some important advantages of SUSY models, such as providing the solutions to the hierarchy problem or Higgs mechanism of electroweak symmetry breaking are lost, but supersymmetry still remains a viable physical basis for inflationary models and dark matter,

as we have demonstrated in this paper on the example of Starobinsky supergravity. To probe this basis, a more detailed analysis of additional cosmological consequences of superhigh energy supersymmetry is needed, while a possibility of implementing the mechanisms of PBH formation can provide a viable test of these consequences. Further analysis of PBH formation and its effects does not only put constraints on the model parameters but can also fix the range of the parameters at which gravitational wave experimental data on massive BH coalescence finds an interpretation in terms of massive PBH clusters.

An implementation of mechanisms of generation of baryon asymmetry in the Universe can convert Starobinsky supergravity into the complete physical basis for the modern cosmology. A possibility of nonhomogeneous baryosynthesis and formation of antimatter domains [3,6,49,76,77,79–83,88–92] can add the set of probes of supersymmetry at super high energy scales, and, perhaps, lead to their experimental evidence in searches for antinuclei in cosmic rays.

We conclude that even if the modern cosmological paradigm implies physics of Starobinsky supergravity at very high energy scales, the methods of cosmoparticle physics can provide a nontrivial set of its indirect cosmological probes.

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