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Scaloron the mighty: Producing dark matter and baryon asymmetry at reheating

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ABSTRACT

In R^2 -inflation scalaron slow roll is responsible for the inflationary stage, while its oscillations reheat the Universe. We find that the same scalaron decays induced by gravity can also provide the dark matter production and leptogenesis. With R^2 -term and three Majorana fermions added to the Standard Model, we arrive at the phenomenologically complete theory capable of simultaneously explaining neutrino oscillations, inflation, reheating, dark matter and baryon asymmetry of the Universe. Besides the seesaw mechanism in neutrino sector, we use only gravity, which solves all the problems by exploiting scalaron.

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1. Introduction and summary

The first inflationary model widely discussed in literature, dubbed R^2 -inflation [1], works in a very economic way, exploiting one and the same interaction—gravity—to accomplish both inflation and subsequent reheating. It is tempting to exploit gravity somewhat further, addressing two other important issues: dark matter production and generation of the baryon asymmetry of the Universe.

To this end we consider in this work the universal mechanism of particle production operating in R^2 -inflationary model: scalaron decay. We have found that free scalars heavier than 10 keV and fermions heavier than 10^7 GeV are forbidden in this scenario, as they would overclose the Universe. The fermion of 10^7 GeV is a viable dark matter candidate in this model, while light scalars could contribute to the *hot* dark matter component at best.

We have also found that with two additional right-handed sterile neutrinos one can explain the neutrino oscillations (via standard seesaw mechanism) and baryon asymmetry of the Universe (via standard non-thermal leptogenesis). Curiously, the amount of baryon asymmetry available in the model is strongly constrained from above, so that the observed amount is only one order of magnitude below the model upper limit.

Both the dark matter fermions and sterile neutrinos are produced in post-inflationary Universe in scalaron decays. This sug-

gests that SM, supplemented with R^2 -term and three right-handed sterile neutrinos, forms a kind of *naturally complete* theory. This theory explains different phenomena beyond the SM—inflation, reheating, dark matter, baryon asymmetry of the Universe—involving *one and the same mechanism* based on the peculiarities of scalaron interactions with itself and other fields. These phenomena together with neutrino oscillations are the main observational facts pointing at incompleteness of the SM. All of them can be explained within the proposed model.

2. Gravitational production of dark matter in R^2 -inflation

We start with the following Lagrangian in the Jordan frame:

$$S^{JF} = -\frac{M_P^2}{2} \int \sqrt{-g} d^4x \left(R - \frac{R^2}{6\mu^2} \right) + S_{matter}^{JF}, \quad (1)$$

where we use the reduced Planck mass M_P related to the Planck mass M_{Pl} as $M_P = M_{Pl}/\sqrt{8\pi} = 2.4 \times 10^{18}$ GeV; S_{matter}^{JF} includes the action of the Standard Model and other new fields. In particular, for free scalar φ and Dirac fermion ψ one has¹

$$S_{\varphi}^{JF} = \int \sqrt{-g} d^4x \left(\frac{1}{2} g^{\mu\nu} \partial_{\mu} \varphi \partial_{\nu} \varphi - \frac{1}{2} m_{\varphi}^2 \varphi^2 \right), \quad (2)$$

$$S_{\psi}^{JF} = \int \sqrt{-g} d^4x (i \bar{\psi} \hat{D} \psi - m_{\psi} \bar{\psi} \psi), \quad (3)$$

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¹ These fields are free in the sense of particle physics: in both the Jordan (Eqs. (2) and (3)) and Einstein frames (Eqs. (5) and (6)), they are free at $M_P \rightarrow \infty$.

where \hat{D} is the covariantly generalized Dirac operator, see e.g. [4]. It is convenient to go to the Einstein frame by the conformal transformation

$$g_{\mu\nu} \rightarrow \tilde{g}_{\mu\nu} = \chi g_{\mu\nu}, \quad \chi = \exp(\sqrt{2/3} \phi/M_P).$$

Scalar and fermion fields are rescaled then as

$$\begin{aligned} \varphi &\rightarrow \tilde{\varphi} = \chi^{-1/2} \varphi, & \psi &\rightarrow \tilde{\psi} = \chi^{-3/4} \psi, \\ \hat{D} &\rightarrow \hat{\tilde{D}} = \chi^{-1/2} \hat{D}. \end{aligned}$$

In the Einstein frame the gravity action takes the Einstein–Hilbert form, but additional scalar degree of freedom ϕ emerges and couples to all matter fields. Thus the original action (1) transforms into

$$S^{EF} = \int \sqrt{-\tilde{g}} d^4x \left[-\frac{M_P^2}{2} \tilde{R} + \frac{1}{2} \tilde{g}^{\mu\nu} \partial_\mu \phi \partial_\nu \phi - \frac{3\mu^2 M_P^2}{4} \left(1 - \frac{1}{\chi(\phi)}\right)^2 \right] + S_{matter}^{EF}, \quad (4)$$

where S_{matter}^{EF} includes interactions with the field ϕ . In particular, now the actions (2) and (3) read

$$S_\varphi^{EF} = \int \sqrt{-\tilde{g}} d^4x \left(\frac{1}{2} \tilde{g}^{\mu\nu} \partial_\mu \tilde{\varphi} \partial_\nu \tilde{\varphi} - \frac{1}{2\chi} m_\varphi^2 \tilde{\varphi}^2 + \frac{\tilde{\varphi}^2}{12M_P^2} \tilde{g}^{\mu\nu} \partial_\mu \phi \partial_\nu \phi + \frac{\tilde{\varphi}}{\sqrt{6}M_P} \tilde{g}_{\mu\nu} \partial_\mu \tilde{\varphi} \partial_\nu \phi \right), \quad (5)$$

$$S_\psi^{EF} = \int \sqrt{-\tilde{g}} d^4x \left(i \tilde{\psi} \hat{\tilde{D}} \tilde{\psi} - \frac{m_\psi}{\sqrt{\chi}} \tilde{\psi} \tilde{\psi} \right). \quad (6)$$

At small values of ϕ this new degree of freedom decouples from all other fields and both frames become identical, as in this limit R is also small.

Cosmology of the homogeneous and isotropic Universe described by the action (1) has inflationary stage [1] at large values of R . In the Einstein frame this stage is realized as large-field inflation in the slow-roll regime taking place at super-Planckian values of the field ϕ which serves as the inflaton. The equivalent scalar mode in the action (1) was named scalaron [1] and we use both names in what follows.

Inflation gives rise to primordial scalar perturbations, whose amplitude normalization to the observed CMB anisotropy and large-scale structure yields the estimate [2]

$$\mu = 1.3 \times 10^{-5} M_P.$$

The spectral index of the scalar perturbations and parameters of the generated tensor perturbations are consistent with observational constraints [3].

When the slow-roll conditions get violated, inflation terminates and the inflaton ϕ starts to oscillate rapidly, with frequency equal to scalaron mass μ . This drives the Universe expansion like at matter-dominated stage. The intermediate stage naturally ends up with inflaton decays into ordinary particles due to *universal* interactions as in (5), (6). In particular, scalaron decay rates into a pair of sufficiently light scalars and into a pair of sufficiently light fermions are²

$$\Gamma_{\phi \rightarrow \varphi\varphi} = \frac{\mu^3}{192\pi M_P^2}, \quad (7)$$

$$\Gamma_{\phi \rightarrow \tilde{\psi}\tilde{\psi}} = \frac{\mu m_\psi^2}{48\pi M_P^2}, \quad (8)$$

respectively. In this *universal* way the energy flows from the inflaton to ordinary particles. There is no any amplification of the energy drain to the SM particles due to coherent effects (like ones in Ref. [6]), since the produced particles interact strongly enough.

Formulas (7), (8) generally mean that scalar particles play the major role in the scalaron decay process. The dominant contribution to the decay rate comes from coupling to kinetic term in Eq. (5). There is no similar term in the fermion case, see Eq. (6), because of the conformal invariance³ of the fermion kinetic term. The energy transfers to the relativistic scalars by the time $t_{reh} \gtrsim 1/\sum_s \Gamma_{\phi \rightarrow \varphi\varphi}$. We define the reheating temperature T_{reh} as an effective temperature of produced relativistic matter at the time of equality between inflaton and matter energy densities. Then in the absence of coherent effects mentioned above, for N_s scalar components contributing to the inflaton decay the effective temperature of produced relativistic matter is

$$T_{reh} \approx 4.5 \times 10^{-2} \times g_*^{-1/4} \cdot \left(\frac{N_s \mu^3}{M_P} \right)^{1/2}, \quad (9)$$

where g_* denotes the number of relativistic species. In case of the Standard Model with $g_* = 106.75$ and $N_s = 4$, Eq. (9) gives numerically

$$T_{reh} \approx 3.1 \times 10^9 \text{ GeV}, \quad (10)$$

which is well in the region where the gauge interactions of the Standard Model are in equilibrium. So, the value (10) is the *maximum* temperature of the primordial plasma or the *reheating* temperature, indeed.

With low reheating temperature (10) the post-inflationary matter-dominated stage lasts long enough, so that the subhorizon scalar perturbations growing at this stage become nonlinear (for a recent study see e.g. [7]). One expects that scalarons then start to form self-gravitating clumps of linear sizes much smaller than horizon, so the Universe still expands as at matter domination. The scalaron overdensity in clumps can be estimated similar to analysis of dark matter halos and clumps in the late-time Universe, see e.g. [8]. The scalaron density in clumps is not high enough to initiate scalaron scatterings due to self-interaction. Therefore, relativistic (SM) particles production can be described as decays of (non-relativistic) scalarons and the estimate of reheating temperature (10) remains intact.

The decay rate formulas (7), (8) show that the scalaron produces ordinary particles in a universal way, so that all scalars and fermions of the theory will eventually populate the expanding Universe. With all non-gravitational interactions between ordinary particles switched off, the resulting abundances would be determined mostly by spin and, for fermions, by mass of the particles.

Now let us consider a new field which is either free in the Jordan frame or couples to other fields very weakly, so that these new particles never equilibrate in the primordial plasma.⁴ If stable at cosmological time scale, this particle is a good candidate to be dark matter.

² These estimates are in agreement with similar ones in the Jordan frame originally obtained in [1,5], while (7) and (8) are two times larger and four times smaller, respectively, than the estimates in Ref. [2].

³ This implies quite a similar situation for a scalar *conformally* coupled to gravity in the Jordan frame. We do not consider this case, but note that the coherent processes at preheating stage we alluded to above may be relevant then.

⁴ In particular, the Planck-scale suppressed nonrenormalizable couplings are allowed.

Given the discussion above, the particle mass is the only free parameter and hence its value is fixed by the requirement of comprising all dark matter whose relative contribution to the present energy density ρ_c is [3] $\Omega_{DM} \approx 0.223$. The mass of dark matter particles m_{DM} and their number density at present $n_{DM,0}$ are related to ρ_c and Ω_{DM} as follows

$$m_{DM} = \frac{\Omega_{DM} \rho_c}{s_0} \frac{s_0}{n_{DM,0}}, \quad (11)$$

where we introduced the present entropy density s_0 .

Let us estimate the entropy-to-dark-matter ratio s/n_{DM} by the time of reheating. The dark matter production after inflation can be described as the decay of non-relativistic scalarons of mass m_ϕ whose number density n_ϕ evolves with scale factor $a = a(t)$ at $t \gtrsim t_{reh}$ as

$$n_\phi(a) = \frac{\mu}{2} \phi_{reh}^2 \left(\frac{a_{reh}}{a} \right)^3.$$

Here ϕ_{reh} , a_{reh} refer to the values of inflaton field and scale factor at reheating. We again neglect any coherent effects related to the new fields and also treat this new decay channel of scalaron as subdominant one, so the reheating temperature is still given by Eq. (10). Both points are discussed below in due course: we will see that the obtained results justify our choice. Assuming two-body decays of the inflaton to dark matter particles with decay rate $\Gamma_{\phi \rightarrow DM}$ one writes down the Boltzmann equation for the dark matter density,

$$\frac{d}{dt} (n_{DM} a^3) = 2n_\phi(a) \Gamma_{\phi \rightarrow DM} a^3.$$

This equation has a solution

$$n_{DM}(t_{reh}) = \frac{\rho_{reh}}{\mu} \Gamma_{\phi \rightarrow DM} t_{reh},$$

where ρ_{reh} is the total energy density at the time of reheating, when the corresponding contributions of inflaton and relativistic species coincide. The energy density is related to the Hubble parameter at reheating H_{reh} by

$$\rho_{reh} = 2g_* \frac{\pi^2}{30} T_{reh}^4 = 3M_P^2 H_{reh}^2.$$

Taking for the numerical estimate of the reheating time

$$t_{reh} \simeq \frac{1}{\sqrt{3} H_{reh}},$$

we finally obtain for the entropy-to-dark-matter ratio at reheating

$$\frac{s}{n_{DM}}(T_{reh}) = \frac{2\pi \sqrt{g_*}}{3\sqrt{15}} \frac{\mu}{\Gamma_{\phi \rightarrow DM}} \frac{T_{reh}}{M_P}. \quad (12)$$

Since this ratio remains intact at the further hot stages, we have the estimate

$$\frac{s}{n_{DM}} \simeq \frac{s_0}{n_{DM,0}}.$$

Eqs. (12), (11) imply the following equation for the mass of dark matter particles universally produced by the scalaron decay,

$$m_{DM} = \frac{\Omega_{DM} \rho_c}{s_0} \frac{2\pi \sqrt{g_*}}{3\sqrt{15}} \frac{\mu}{\Gamma_{\phi \rightarrow DM}} \frac{T_{reh}}{M_P}. \quad (13)$$

Then in the cases of scalar (7) and fermion (8) one obtains the numerical estimates

$$m_\phi \approx (6.9 \text{ keV}) \left(\frac{1.3 \times 10^{-5} M_P}{\mu} \right)^{1/2} \left(\frac{N_s}{4} \right)^{1/2} \times \left(\frac{g_*}{106.75} \right)^{1/4} \left(\frac{\Omega_{DM}}{0.223} \right) \left(\frac{\rho_c}{0.52 \times 10^{-5} \text{ GeV/cm}^3} \right), \quad (14)$$

$$m_\psi \approx (1.2 \times 10^7 \text{ GeV}) \left(\frac{\mu}{1.3 \times 10^{-5} M_P} \right)^{1/2} \left(\frac{N_s}{4} \right)^{1/6} \times \left(\frac{g_*}{106.75} \right)^{1/12} \left(\frac{\Omega_{DM}}{0.223} \right)^{1/3} \times \left(\frac{\rho_c}{0.52 \times 10^{-5} \text{ GeV/cm}^3} \right)^{1/3}. \quad (15)$$

Here we have corrected numerical coefficients (by tens of percent) to match the results obtained with numerical integration of the Boltzmann equations which consistently describe both the energy transfer from scalaron to all relativistic scalars and the dark matter production.

The dark matter particles are produced highly relativistic with 3-momenta $\sim \mu/2$, which exceed T_{reh} and hence momenta of particles in the plasma by about four orders of magnitude,

$$\mu \sim 10^4 \times T_{reh}.$$

So, the free scalars of 10 keV mass would contribute to *hot dark matter* component only. This component is certainly subdominant, hence free scalars cannot solve the dark matter problem in the case of R^2 -inflation. The viable choice of dark matter is free fermions of mass $m_{DM} \approx 10^7$ GeV. These particles would form *cold dark matter*.

We end up the discussion of universal dark matter in R^2 -inflationary model with several remarks. First, for the considered case of free boson particles coherent effects [6] during scalaron oscillations could change the obtained results. However, it only makes worse the situation with scalar dark matter, and we do not consider these processes.

Second, if at a later stage of the Universe expansion the entropy gets produced (say, due to decays of nonrelativistic particles out of equilibrium, etc.), the estimates (14), (15) have to be corrected to account for the corresponding dilution factor $r_s = s_{new}/s_{old}$. Thus, r.h.s. of (14) has to be divided by r_s , while r.h.s. of (15) has to be divided by $r_s^{1/3}$. In realistic models with moderate (if any) entropy production this does not save scalar dark matter, and only mildly changes the prediction for the fermion dark matter mass (15).

Third, in the extensions of particle physics with additional scalars coupled to the SM fields, both reheating temperature (9) and would-be scalar dark matter mass increase. With very large number of new scalars this makes the scalar dark matter viable.

Fourth, if there are other sources of dark matter in the model—other stable particles or other mechanisms of out-of-equilibrium production of heavy fermions—the estimate (15) implies the *upper limit* on the mass of universal dark matter in R^2 -inflationary model.

Fifth, in case of the Majorana fermions the mass is larger by a factor of $2^{1/3}$ as compared to the Dirac case (15), since the decay rate of scalaron into the Majorana particles is two times smaller.

Sixth, the estimates (14), (15) mean that *heavier free particles are forbidden* in models with R^2 -inflation, otherwise the Universe would be overclosed. This conclusion is true for the particles lighter than the scalaron, when these estimates of particle production rates are applicable. Clearly, production rate of heavier particles by means of scalaron oscillations is suppressed as compared to (7), (8). The heavier the particles, the stronger the suppression. Hence, free particles of mass $m \gg \mu$ are not forbidden in R^2 -inflation. Moreover, with strongly (most probably exponentially)

tuned value of mass $m \sim \mu$ they can form viable dark matter. We do not study this situation here.

3. Leptogenesis in R^2 -inflation

The study presented in the previous section revealed that heavy free (in the Jordan frame) fermion is a viable dark matter candidate in R^2 -inflationary model. These fermions are naturally produced in the post-inflationary Universe by scalaron decays, so the right amount of dark matter is achieved with the mass of about 10^7 GeV (15). One can further ask whether it is possible to make use of this universal production mechanism to unravel another cosmological problem which SM fails to solve: the baryon asymmetry of the Universe. The answer is positive, and we illustrate it with example of nonthermal leptogenesis via decays of heavy sterile neutrinos [9] universally produced by scalaron decays. This particular example is strongly motivated by observed oscillations of active neutrinos, phenomena which also lack an explanation within SM. Heavy sterile neutrinos give masses to active neutrinos via seesaw mechanism. This mechanism is responsible for the hierarchy between the neutrino masses and electroweak scale. In this Section we show that both the correct neutrino masses and successful leptogenesis can be realized in the R^2 -inflationary model.

The modification of SM we consider consists of two new Majorana fermions N_I , $I = 1, 2$, which are right singlets with respect to the SM gauge group. The most general renormalizable Lagrangian for these fermions is

$$\mathcal{L} = i\bar{N}_I \gamma^\mu \partial_\mu N_I - y_{\alpha I} \bar{L}_\alpha N_I \tilde{\Phi} - \frac{M_I}{2} \bar{N}_I^c N_I + \text{h.c.}, \quad (16)$$

where $y_{\alpha I}$ are new Yukawa couplings, Φ is the SM Higgs doublet and $\tilde{\Phi} = \epsilon \Phi^*$ with ϵ being antisymmetric 2×2 matrix.

When electroweak symmetry breaks down, Φ acquires vacuum expectation value, $\Phi^T = (0, v/\sqrt{2})$, $v = 246$ GeV. Then (16) yields active-sterile mixing in the neutrino mass matrix. Assuming $yv \ll M_I$, this mixing gives for the active neutrino masses:

$$m_{\nu\alpha\beta} = - \sum_I y_{\alpha I} \frac{v^2}{2M_I} y_{\beta I}. \quad (17)$$

Hence, $m_\nu \ll v$, that is the seesaw mechanism.⁵ The formula (17) together with atmospheric neutrino mass splitting, $\Delta m_{atm}^2 \simeq 3 \times 10^{-3} \text{ eV}^2$ [10], implies an order-of-magnitude upper limit on the lightest sterile neutrino mass:

$$M_{lightest} \lesssim \frac{v^2}{2\sqrt{\Delta m_{atm}^2}} \simeq 5 \times 10^{14} \text{ GeV}. \quad (18)$$

The smaller the Yukawa couplings $y_{\alpha I}$ and the heavier the active neutrinos, the smaller $M_{lightest}$, see Eq. (17).

Lagrangian (16) added to SM explicitly violates the lepton number. The source of CP-violation, which is one of the Sakharov's conditions [11] for successful baryo- and leptogenesis, is the complex Yukawa couplings $y_{\alpha I}$. The third Sakharov's condition, departure from thermal equilibrium for lepton number violating processes, is achieved due to decay of non-relativistic sterile neutrinos, which is evidently an out-of-equilibrium process.

Let us discuss the production of these heavy sterile neutrinos in the post-inflationary Universe with oscillating scalaron. Neutrinos are produced via Yukawa-type coupling to scalaron, see Eq. (6). The production is effective only for neutrinos lighter than scalaron.

⁵ Note that with only two seesaw sterile neutrinos we end up with one out of three active neutrino being massless.

Indeed, the amplitude of scalaron oscillation never exceeds M_p . At the same time the scalaron coupling to a fermion is proportional to the fermion mass, so the latter remains almost intact. Therefore, an enhancement of heavy fermion production by Yukawa source due to periodic decrease of fermion effective mass [12] does not work here.

Thus, we consider the models with sterile neutrinos lighter than scalaron, $M_I < m_\phi$. By the time of reheating, $t \simeq t_{reh}$, the ratio of number density of produced in scalaron decays heavy neutrinos to entropy is given by Eq. (12). Substituting relevant numbers in (12) and (8) (recall that for Majorana fermions the decay rate is two times lower than for the Dirac fermions) one obtains

$$\frac{n_{N_I}(T_{reh})}{s} = 2.9 \times 10^{-6} \times \left(\frac{M_I}{5 \times 10^{12} \text{ GeV}} \right)^2, \quad (19)$$

where we have included a correction of about tens of percent in order to match the solution of the relevant Boltzmann equations (see comment below Eqs. (14), (15)).

These sterile neutrinos, however, decay quite rapidly due to the same Yukawa couplings in (16) responsible for the seesaw mechanism. Their total decay rates are

$$\Gamma_{N_I} = \frac{M_I}{8\pi} \sum_\alpha |y_{\alpha I}|^2.$$

Disregarding any hierarchy in the neutrino Yukawa sector one takes for the order-of-magnitude numerical estimate of Yukawa couplings relation (17) with $m_\nu^2 \sim \Delta m_{atm}^2$ and finds for the lightest sterile seesaw neutrino

$$\Gamma_{N_I} \sim \frac{\sqrt{\Delta m_{atm}^2} M_I^2}{4\pi v^2}.$$

Thus, sterile neutrino of mass above 10^{10} GeV decays before reheating.⁶ This is a very process which generates lepton asymmetry in the early Universe. To simplify formulas we assume further, that only the lightest neutrino, N_1 , contributes and $M_1 \ll M_2$. Then the value of lepton asymmetry $\Delta_L = n_L/s$ is

$$\Delta_L = \delta_L \cdot \frac{n_{N_1}}{s}, \quad (20)$$

with [9] (for reviews see e.g. [13])

$$\begin{aligned} \delta_L &= - \frac{3M_1}{8\pi v^2} \cdot \frac{1}{\sum_{\alpha, I} |y_{\alpha I}|^2} \sum_{\alpha\beta} \text{Im} \left(y_{\alpha 1} y_{\beta 1} y_{\alpha 2}^* \frac{v^2}{2M_2} y_{\beta 2}^* \right) \\ &\sim \frac{3M_1 \sqrt{\Delta m_{atm}^2}}{8\pi v^2}, \end{aligned} \quad (21)$$

where we have used the order-of-magnitude estimate (17) with $m_\nu^2 \sim \Delta m_{atm}^2$. Finally, from (20), (21) and (19) we obtain

$$\Delta_L \sim 1.5 \times 10^{-9} \times \left(\frac{M_1}{5 \times 10^{12} \text{ GeV}} \right)^3. \quad (22)$$

The asymmetry is generated mostly at late times just before reheating, when the ratio of sterile neutrino (if stable) density to entropy becomes maximum, see Eq. (12). The sterile neutrinos are quite heavy so the lepton asymmetry is not washed out by inverse decay processes, which are out of equilibrium. Likewise, another important lepton violating process, scatterings $\Phi L \leftrightarrow \Phi \bar{L}$ through

⁶ Eqs. (19) and (10) mean that sterile neutrino decays contribute not much to the Universe reheating.

virtual sterile neutrinos (with cross section $\sigma \sim \frac{m_\nu^2}{8\pi v^4}$) are out-of-equilibrium at the epoch of interest. Thus, lepton asymmetry (22) transfers later on into baryon asymmetry by sphaleron processes [14],

$$\Delta_B \simeq \Delta_L/3 \sim 0.5 \times 10^{-9}.$$

Given the present baryon asymmetry of the Universe $\Delta_{B,0} = 0.88 \times 10^{-10}$ [3] and the estimates above, we conclude that heavy seesaw sterile neutrinos of mass 10^{12} – 10^{13} GeV, produced in scalaron decays, in R^2 -inflationary model can be responsible for the baryon asymmetry of the Universe. Since the sterile neutrino are lighter than scalaron, the available room is not large, that one could interpret as *an order-of-magnitude prediction of the baryon asymmetry*, which is otherwise (and usually) appears as an accidental number.

The estimates above are an order-of-magnitude only. In the particular model under discussion with only two seesaw sterile neutrinos one can relate all the Yukawa couplings to the measured active neutrino mass differences and mixing angles, so that only few model parameters remain free. In this way one can refine our estimates. Likewise the numerical estimates may be changed in models with a particular hierarchy of Yukawa couplings in the neutrino sector.

Note in passing that second, heavier sterile neutrino N_2 if not too heavy can also contribute to lepton asymmetry generation, its decay products do not equilibrate. There is also some contribution from sterile neutrino decays at earlier times, when they are relativistic. The inverse processes are rare because the particle density is small.

4. Discussion

To conclude, we have shown that in R^2 -inflation the scalaron decay responsible for reheating can also take care of dark matter production and generation of the baryon asymmetry of the Universe. The latter occurs through nonthermal leptogenesis by introducing seesaw neutrinos. Thus, neutrino oscillations are also explained in this model, which pretends to be *phenomenologically self-contained*.

We end up with several comments. First, at relatively low energy scale, the model we discussed can be further modified in gravity sector (attempting to explain the late-time accelerated expansion of the Universe, as e.g. in Refs. [15,16]) or in particle physics sector (say, by adding PQ -axion [17] to solve the strong CP-problem). All these and similar modifications can be safely adopted provided the early time cosmology remains intact.

Second, we do not discuss here the gauge hierarchy problem in particle physics. Certainly, R^2 -term does not contribute to this problem: all couplings of scalaron to other particles are suppressed by the Planck scale. With the absence of new scales in *particle physics*, one cannot discard that still unknown complete quantum theory at gravity scale (quantum gravity) is responsible for the cancellation of dangerous quantum corrections (if any) to the Higgs boson mass. In fact, a nontrivial space structure that shows up at Planck scales could also invalidate the strong CP-problem (see e.g. [18]). Thus, quantum gravity could solve all problems, indeed. Otherwise, one can think about other solutions usually considered. In particular, our logic can be adopted for supersymmetric extensions of the SM and models with axion. In all these cases free fermion again serves as dark matter candidate, though it is somewhat heavier than in the case of SM because of larger number N_s of scalars responsible for reheating and larger number of degrees of freedom g_* , see Eq. (15).

Third comment, which is related to the second one. The heavy sterile neutrinos (of mass M_N) we used contribute to gauge hierarchy problem making it worse as compared to the SM. Indeed, their coupling to the SM Higgs boson gives rise to both divergent and finite, of order M_N , contributions to the Higgs boson mass. This worsening can be avoided if sterile neutrinos coupled to the Higgs boson have masses of order electroweak scale or smaller, as e.g. in ν MSM model [19]. Indeed, heavy sterile neutrinos can be replaced by the light ones in R^2 -model we discuss. In that case, however, sterile neutrinos, required for leptogenesis, should be produced not by the scalaron decay (which, as we saw, requires large mass to be efficient enough), but via neutrino oscillations in the early Universe [20]. The R^2 -model with one heavy fermion (free in the Jordan frame) and two light seesaw sterile neutrinos can be *both phenomenologically and theoretically self-contained* upto the quantum gravity effects. At the same time, this modification has also one more advantage as compared to what we discuss in the main text: sterile neutrino sector here can be directly tested [21] in particle physics experiments.

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