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Light inflaton hunter's guide

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Abstract: We study the phenomenology of a realistic version of the chaotic inflationary model, which can be fully and directly explored in particle physics experiments. The inflaton mixes with the Standard Model Higgs boson via the scalar potential, and no additional scales above the electroweak scale are present in the model. The inflaton-to-Higgs coupling is responsible for both reheating in the Early Universe and the inflaton production in particle collisions. We find the allowed range of the light inflaton mass, $270\text{ MeV} \lesssim m_\chi \lesssim 1.8\,\text{GeV}$, and discuss the ways to find the inflaton. The most promising are two-body kaon and B-meson decays with branching ratios of orders 10^{-9} and 10^{-6}, respectively. The inflaton is unstable with the lifetime 10^{-9}–10^{-10} s. The inflaton decays can be searched for in a beam-target experiment, where, depending on the inflaton mass, from several billions to several tenths of millions inflatons can be produced per year with modern high-intensity beams.

Keywords: Cosmology of Theories beyond the SM, Rare Decays

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Contents

1 Introduction

In this paper we present an example of how (low energy) particle physics experiments can directly probe the inflaton sector (whose dynamics is important at high energies in the very Early Universe).

The common assumption about the inflaton sector is that it is completely decoupled from the Standard Model (SM) at energies much lower than the inflationary scale. This assumption appears quite natural, since the slow-roll conditions generally permit only tiny coupling of the inflaton to any other fields including itself (see e.g. [1] for a review). This implies within the perturbative approach that by integrating out the inflaton sector one obtains at low energies some non-renormalisable operators strongly suppressed by both tiny couplings and high inflationary scale. These assumptions prevent any direct (laboratory) investigation of the inflationary mechanism.

The situation is quite different if the inflaton sector contains only light fields. In this case even weak inflaton coupling to the SM particles can lead to observable signatures in laboratory experiments.

We will concentrate here on the model, which is a particular version of the simple chaotic inflation with quartic potential and with inflaton field coupled to the SM Higgs boson via a renormalisable operator. We confine ourselves to the case where the Higgs scalar potential is scale-free at the tree level, so that its vacuum expectation value is proportional to that of the inflaton field and Higgs boson mixes with inflaton. In this model, the inflaton χ is found to be light,[1] 270 MeV $\lesssim m_\chi \lesssim$ 10 GeV, where the lower limit actually comes

[1] In fact, there is another window for the inflaton mass, which is also below 1 TeV, see eqs. (2.9), (2.10).

from the searches for decays $K \to \pi +$ nothing and from the search for axion-like particles in CHARM experiment, and the upper limit is related to the requirement of a sufficient reheating after inflation. The most promising here are searches of inflaton in decays of K- and B-mesons and searches for inflaton decays in beam-target experiments.

We present the estimates of meson decay branching ratios into the inflaton, and give the estimate for the inflaton production rate in beam-target experiment with beam parameters of the T2K, NuMi, CNGS and NuTeV experiments [2–5]. The interesting branchings start from 10^{-6} for the decay $B \to K + \chi$ and 10^{-9} for the decay $K \to \pi + \chi$ (see section 4). The production rate in the beam-target experiments is of several millions per year (at $m_\chi \sim 500\,\text{MeV}$) to thousands (at $m_\chi \sim 5\,\text{GeV}$). Models with larger masses are hard to explore. However, these larger masses correspond to quite heavy SM Higgs boson, $m_h \sim 350$–$700\,\text{GeV}$. For the described model being consistent up to the inflationary scales the Higgs mass should be $m_h < 190\,\text{GeV}$, and the inflaton mass is below $m_\chi = 1.8\,\text{GeV}$ (see eq. (2.8)). Hence we conclude that in the model considered here the inflaton sector can be fully explored in the particle physics experiments.

In this respect it is worth mentioning, that hypothetical light bosons, (very) weakly coupled to the SM fields, are present in various extensions of the SM (for particular examples see [6–9]). Phenomenology of these particles attracted some attention, and in ref. [10] the list of relevant experiments is given in the section about searches for axions and other very light bosons. These experiments are relevant for the described model (see the bounds obtained in this paper). At the same time the theoretical development for our case of a light flavour blind scalar dates back to the time, when very light SM Higgs boson was still allowed experimentally. Quite a few improvements were done since that time, and we mostly used in the current work the results obtained in refs. [11–15].

The rest of the paper is organised as follows. In section 2 we describe the model and outline the viable region in the parameter space, section 3 is devoted to inflaton decays, while meson decays to inflaton are considered in section 4. There we obtain lower limits on the inflaton mass from existing experimental results. In section 5 we study the inflaton production in pp-collisions and give production rates for operating high-intensity and high-energy beams at JPARC, Fermilab and CERN. Limits on the inflaton mass from the results of the CHARM experiment are obtained in section 6. There we also give predictions for the number of inflatons produced per one year of operation of T2K, NuMi, CNGS and NuTeV beams in a beam-target setup. Section 7 contains conclusions. In appendix A we discuss consequences of obtained results for a implementation of our model within the νMSM [16], the extension of the SM with three sterile right-handed neutrinos, so that the inflaton vacuum expectation value provides sterile neutrino masses. This model is capable of explaining neutrino oscillations, dark matter and baryon asymmetry of the Universe and, equipped with inflaton sector, provides an example of a full realistic model of particle physics. Thus the inflation can be directly tested also in fully realistic extensions of the SM.

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2 The model

We consider the extension of the SM model with an inflaton field introduced in [16]. The Lagrangian of this extended model reads

$$\begin{aligned}\mathcal{L}_{XSM} &= \mathcal{L}_{\rm SM} + \mathcal{L}_{XN}\;,\\ \mathcal{L}_{XN} &= \frac{1}{2}\partial_\mu X \partial^\mu X + \frac{1}{2}m_X^2 X^2 - \frac{\beta}{4}X^4 - \lambda\left(H^\dagger H - \frac{\alpha}{\lambda}X^2\right)^2\;, \end{aligned} \tag{2.1}$$

where $\mathcal{L}_{\rm SM}$ is the SM Lagrangian without the Higgs potential, while the latter gets modified in accordance with (2.1), X is a new neutral scalar field and H is the Higgs doublet. Note, that it is supposed that the only scale violating term at tree level is the mass term with the negative squared mass $-m_X^2$ for the extra scalar X. This particular choice is sufficient to demonstrate the main statement: possibility to test directly the inflationary models in particle physics experiments. At the same time, the algebra below is simpler as compared to those in the case of general inflaton-Higgs scalar potential, since the number of parameters is smaller.

The SM-like vacuum of the scalar potential in eq. (2.1) is

$$\begin{aligned} &\langle H\rangle = \frac{v}{\sqrt{2}}\;, \qquad\qquad \langle X\rangle = \sqrt{\frac{\lambda}{2\alpha}}\,v = \frac{m_\chi}{\sqrt{2\beta}}\;,\\ &\text{with } v = \sqrt{\frac{2\alpha}{\beta\lambda}}\,m_X = 246\text{ GeV}\;. \end{aligned} \tag{2.2}$$

The small excitations about this vacuum have the masses

$$m_h = \sqrt{2\,\lambda}\,v\;, \quad m_\chi = m_h\sqrt{\frac{\beta}{2\alpha}}\;, \tag{2.3}$$

and are rotated with respect to the gauge basis $(\sqrt{2}H - v, X)$ by the angle

$$\theta = \sqrt{\frac{2\alpha}{\lambda}} = \frac{\sqrt{2\beta}\,v}{m_\chi} \tag{2.4}$$

to the leading order in θ and m_χ/m_h.

The inflation in the model is supposed to be driven by a flat potential along the direction

$$H^\dagger H \cong \frac{\alpha}{\lambda}X^2\;. \tag{2.5}$$

The quartic coupling dominates the potential during inflationary and reheating stages. Thus, after inflation the Universe expands as at radiation dominated stage, so the number of required e-foldings is about $N_e \simeq 62$. The normalisation of the matter power spectrum [17] generated during inflation fixes[2] the quartic coupling constant as $\beta \approx \beta_0$, where

$$\beta_0 = 1.5\times 10^{-13}\;. \tag{2.6}$$

[2] The estimate (2.6) differs from the other values, presented in [16] and in [18] due to different values of N_e used there.

In this respect it is worth noting, first, that inflation in this model happens at $\chi \gtrsim M_{\rm Pl}$, where gravity-induced corrections are expected to be large. We will suppose below, that same (yet unknown) mechanism, which guarantees the flatness of scalar potential during inflation, operates at $\chi \lesssim M_{\rm Pl}$ as well. Thus, the Plank scale physics allows for considering the same scalar potential (2.1) at high energies with account of perturbative quantum corrections only due to the gauge and Yukawa couplings. These latter are discussed in due course. Second, for the quartic inflation there is a tension [17] between predicted tensor-to-scalar amplitudes ratio and fits to cosmological data. Though the limits [17] are not dramatic yet, weak non-minimal coupling to Ricci scalar, $\xi X^2 R/2$, with $\xi \sim 10^{-3}$ makes the model fully consistent [17, 19, 20] with all current cosmological observations. Switching on ξ results in a larger value of the quartic coupling constant, $\beta_0 \le \beta \lesssim 2\beta_0$ for $0 \le \xi \lesssim 10^{-3}$. For our study of inflaton phenomenology at low energies non-minimal coupling to gravity is irrelevant, but in further estimates we account for the uncertainty in the value of β,

$$\beta = (1\text{–}2) \cdot \beta_0 \;, \tag{2.7}$$

associated with its possible impact.

Thus, among four parameters in the Lagrangian (2.1), one, β, is fixed by the amplitudes of primordial perturbations and another combination (2.2) is fixed at the electroweak vacuum by the value of the Fermi constant. Two remaining free parameters determine the SM Higgs boson mass and the inflaton mass. Further constraints on them are discussed below.

The baryon asymmetry of the Universe is unexplained within the framework of the Standard model of particle physics (SM). However, the baryon number is violated at microscopic level in primordial plasma, if sphaleron processes are rapid enough [21]. This phenomena is often exploited by mechanisms generating baryon asymmetry within relevant extensions of the SM. This places a lower bound on the reheating temperature of the Universe at the level somewhat above the electroweak scale, and for definiteness we choose $T_r \gtrsim 150\,\text{GeV}$ (for a review see [22]). In the model (2.1), where both quartic coupling and Higgs-to-inflaton mixing are very weak, but Higgs boson self-coupling is quite strong ($\lambda \gtrsim 0.1$ for $m_h \gtrsim 114\,\text{GeV}$), the energy transfer from the inflaton to the SM particles is extremely inefficient [18]. The stronger is the mixing, the more efficient is the energy transfer and the higher is the reheat temperature in the early Universe. Strong mixing (larger α) implies lighter inflaton, see eq. (2.3), and lower bound $T_r \gtrsim 150\,\text{GeV}$ yields *upper bounds* on the inflaton mass [18]:

$$m_\chi \lesssim 1.5 \cdot \left(\frac{m_h}{150\ \text{GeV}}\right) \cdot \left(\frac{\beta}{1.5\times 10^{-13}}\right)^{1/2}\ \text{GeV for } m_\chi < 2m_h \;, \tag{2.8}$$

and

$$m_\chi \lesssim 460 \cdot \left(\frac{m_h}{150\ \text{GeV}}\right)^{4/3} \cdot \left(\frac{\beta}{1.5\times 10^{-13}}\right)^{1/3}\ \text{GeV for } m_\chi > 2m_h \;. \tag{2.9}$$

In the latter case the *lower bound* on the inflaton mass is

$$m_\chi > 2m_h \gtrsim 228\ \ \text{GeV} \;. \tag{2.10}$$

In the former case the *lower bound* on the inflaton mass follows from the upper limit on the Higgs-inflaton mixing, appearing from the requirement that quantum corrections, originated from this mixing, should not dominate over bare coupling constant β. With the action (2.1) fixed at the electroweak scale the corrections to the inflationary potential $\beta X^4/4$ can be explicitly calculated and have the form

$$\delta V = \frac{m^4(X)}{64\pi^2} \log \frac{m^2(X)}{\mu^2} \; , \tag{2.11}$$

where $m(X)$ is the mass of the contributing particle in the inflaton background field X (taking into account the flat direction (2.5) to obtain the Higgs field background, and μ is the electroweak scale.[3] Then, requiring that in the inflationary region $X \sim M_P$ the corrections to the quartic coupling β are, somewhat arbitrary, smaller than 10%, we get from the contribution of the Higgs boson

$$\alpha \lesssim 10^{-7} \; , \tag{2.12}$$

which precludes large quantum corrections to inflaton quartic coupling driving inflation. Limit (2.12) can be converted using (2.3) into the lower bound on the inflaton mass

$$m_\chi > 120 \left(\frac{m_h}{150 \;\; \text{GeV}} \right) \left(\frac{\beta}{1.5 \times 10^{-13}} \right)^{\frac{1}{2}} \left(\frac{10^{-7}}{\alpha} \right)^{\frac{1}{2}} \; \text{MeV} \; . \tag{2.13}$$

Below this limit one should take into account quantum corrections to the inflationary potential, which may change the value of the inflaton coupling constant at the electroweak scale (2.6), or even spoil the inflationary picture.

The proper renormalisation group enhancement of the analysis should be done, once any experimental evidence for the light inflaton is found. However, in the larger part of the parameter space no significant changes to the described bounds are expected. Limits, similar to (2.12), follow from the requirement of smallness of the SM gauge and Yukawa coupling corrections. As far as all the SM particle masses during inflation are proportional to $\sqrt{\alpha/\lambda}X$ (see (2.5)), all these bounds differ from (2.12) only by ratios of the coupling constants of the form $y_t/\sqrt{\lambda}$, $g/\sqrt{\lambda}$, etc. This changes the lower bound (2.13). Note, that the exact value of this bound is not crucial due to the stronger experimental constraints, obtained in section 6. However, in some regions of parameter space these corrections may lead to much stronger effects — for example, for special ranges of the Higgs boson mass the Higgs self-coupling becomes small at inflationary scales, so $\sqrt{\alpha/\lambda}X$ turns to be large, hence so do the SM perturbative corrections to the inflaton potential.

The weakness of Higgs-to-inflaton mixing (2.12) is responsible for very tiny inflaton interaction with SM particles. This makes searches for heavy inflaton, in the range given by (2.9), (2.10), hopeless in foreseeable future. Indeed, inflaton direct production requires to collect and study enormously large statistics in high energy collisions. In contrast, the opposite case of light inflaton in the range between (2.8) and (2.13) is quite promising, since

[3] Strictly speaking this is the contribution of a scalar particle, while the numeric coefficient changes with the number of spin degrees of freedom, and also changes the sign for fermionic particles.

inflaton production does not require very high energy at a collision point. Then inflaton can be produced in beam-target experiments, where large statistics is achievable. In this paper we consider low energy phenomenology of this light inflaton.

In next sections we will estimate the decay and production rates of the light inflaton. Since inflaton couplings do not depend on the SM Higgs boson mass, its value determines only viable inflaton mass range (2.8), (2.13). Considering for the Higgs boson mass the range 114 GeV $< m_h <$ 700 GeV as still possible, in what follows we study the inflaton low energy phenomenology for its mass in the interval

$$30 \text{ MeV} \lesssim m_\chi \lesssim 10 \text{ GeV} . \tag{2.14}$$

Actually, in the model under consideration the upper limit on the SM Higgs boson mass is lower than 700 GeV, as we use the same scalar sector (2.1) to describe the inflation at high energies. Indeed, if one considers the inflationary model (2.1) as it is, it should be valid (does not become strongly coupled) up to the energy scale $\sqrt{\alpha/\lambda}X \sim 10^{15}$ GeV. As far as the inflaton is very weakly coupled with the Higgs field, this requirement is the same, as for the Standard Model, leading to the bound $m_h \lesssim 190$ GeV (see, e.g. [23]), or

$$m_\chi \lesssim 1.8 \text{ GeV} . \tag{2.15}$$

Also, the fit to the electroweak data points at mass interval $m_h < 285$ GeV [10]. Nevertheless, in extensions of this model the upper limit on the Higgs boson mass may be higher, so we will discuss the whole interval (2.14) to make our study applicable in a more general case.

3 Inflaton decay palette

Light inflaton decays due to the mixing with the SM Higgs boson (2.4). Thus, its branching ratios coincide (taking into account the small mixing angle (2.4)) with those of the light SM Higgs boson, studied in [11]. In what follows we actually update the results of [11] in view of further relevant developments and findings.

Inflaton of the mass below 900 MeV decays mostly into

$$\gamma\gamma,\ e^+e^-,\ \mu^+\mu^-,\ \pi^+\pi^-,\ \pi^0\pi^0 .$$

Mixing (2.4) gives rise to the Yukawa-type inflaton coupling to the SM fermions f,

$$\mathcal{L}_{\chi\bar{f}f} = \theta\,\frac{m_f}{v}\,\chi\bar{f}f = \sqrt{2\beta}\,\frac{m_f}{m_\chi}\chi\bar{f}f , \tag{3.1}$$

where m_f is the fermion mass. Effective inflaton-pion interaction follows from the Higgs boson coupling to the trace of the energy-momentum tensor, (cf. [13, 14])

$$\begin{aligned} \mathcal{L}_{\chi\pi\pi} = {} & 2\kappa\sqrt{2\beta}\cdot\frac{\chi}{m_\chi}\cdot\left(\frac{1}{2}\partial_\mu\pi^0\partial^\mu\pi^0 + \partial_\mu\pi^+\partial^\mu\pi^-\right) \\ & -(3\kappa+1)\sqrt{2\beta}\cdot\frac{\chi}{m_\chi}\cdot m_\pi^2\cdot\left(\frac{1}{2}\pi^0\pi^0 + \pi^+\pi^-\right) , \end{aligned} \tag{3.2}$$

with m_π being the pion mass and

$$\kappa = 2\frac{N_h}{3b} = \frac{2}{9} \, , \tag{3.3}$$

where $N_h = 3$ is the number of heavy flavours, $b = 9$ is the first coefficient in the QCD beta function without heavy quarks. Finally, Higgs-inflaton mixing results in the inflaton coupling to W-boson, which contributes to the triangle one-loop diagram responsible for the inflaton decay to a pair of photons. Similar contributions come from fermion loops, so that inflaton decay to photons can be described by the effective Lagrangian (cf. [12])

$$\mathcal{L}_{\chi\gamma\gamma} \approx \frac{F\alpha}{4\pi}\frac{\sqrt{2\beta}}{m_\chi}\chi F_{\mu\nu}F^{\mu\nu} \, , \tag{3.4}$$

where α is the fine structure constant, F is a sum of loop contributions from W-boson and fermions f with electric charge eq_f, $F = F_W + \sum_{f,\text{colors}} q_f^2 F_f$

$$\begin{aligned} F_W &= 2 + 3y\left[1 + (2-y)\,x^2\right] \, , \\ F_f &= -2y\left[1 + (1-y)\,x^2\right] \, , \end{aligned} \tag{3.5}$$

and $y = 4m^2/m_\chi^2$, with m being the mass of the contributing particle; here also

$$\begin{aligned} x &= \text{Arctan}\frac{1}{\sqrt{y-1}} & &, \text{ for } y > 1 \, , \\ x &= \frac{1}{2}\left(\pi + i\,\log\frac{1+\sqrt{1-y}}{1-\sqrt{1-y}}\right) & &, \text{ for } y < 1 \, . \end{aligned}$$

In fact, for the interesting range of the inflaton mass the fermion contributions almost cancel the W-boson contribution (the latter dominates over a contribution of each single fermion).

The inflaton-to-SM fields couplings presented above yield the inflaton decays to leptons with the rates:

$$\Gamma_{\chi\to\bar{l}l} = \frac{\beta m_l}{4\pi}\frac{m_l}{m_\chi}\left(1 - \frac{4m_f^2}{m_\chi^2}\right)^{3/2} \, , \tag{3.6}$$

the inflaton decays to pions with the rates:

$$\Gamma_{\chi\to\pi^+\pi^-} = 2\Gamma_{\chi\to\pi^0\pi^0} = \frac{\beta m_\chi}{8\pi}\cdot\left(\frac{2}{9} + \frac{11}{9}\frac{m_\pi^2}{m_\chi^2}\right)^2\sqrt{1 - \frac{4m_\pi^2}{m_\chi^2}} \, , \tag{3.7}$$

and the inflaton decay to photons with the rate:

$$\Gamma_{\chi\to\gamma\gamma} = |F|^2\left(\frac{\alpha}{4\pi}\right)^2\frac{\beta\, m_\chi}{8\pi} \, . \tag{3.8}$$

Note, the tree-level estimate (3.7) is not correct far from the pion threshold, where strong final state interactions become important [15]. Thus, in our numerical estimates we follow ref. [15] to improve (3.7).

If inflaton is heavier than 900 MeV, its hadronic decay modes become more complicated. First, each quark flavour contributes to the inflaton decay if the inflaton mass

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exceeds the double mass of the lightest hadron of the corresponding flavour.[4] Strange quark starts to contribute for the inflaton mass $m_\chi > 2m_K$, where m_K is the K-meson mass. Charm quarks contribute, if $m_\chi > 2m_D$, where m_D is the D-meson mass. Lightest flavour states are mesons and close to the thresholds the inflaton decays are described by effective interactions similar to (3.2). Farther from the thresholds final state interactions become important and we follow ref. [15] to estimate the decay branching rate to pions and kaons. This approach is valid until other flavour resonances enter the game or while multimeson final states are negligible. For pions this approach becomes unjustified in models with inflaton mass above approximately 1 GeV. For mesons of heavier flavours the approach fails not far from the corresponding thresholds, where produced mesons are not relativistic but many new hadrons enter the game. Far above a quark threshold inflaton decay can be described as decay into a pair of quarks due to coupling (3.1), which subsequently hadronise. To estimate the decay rate, QCD-corrections should be accounted for, which are the same as those for the SM Higgs boson of mass m_χ.

Second, if the inflaton mass is close to the mass of some narrow hadronic scalar resonance (e.g., oniums in $c\bar{c}$ and $b\bar{b}$ systems), then interference with such a resonance changes significantly inflaton hadronic decay rates in a way quite similar to what was expected for the SM Higgs boson if its mass would be close to the $b\bar{b}$ threshold [24]. We do not consider here this case, though the calculations are straightforward.

Third, for the inflaton mass above 1 GeV the decay into mesons actively proceeds via the inflaton-gluon-gluon coupling, which due to the quark one-loop triangle diagram similar to that contributing to the inflaton decays to photons. The effective Lagrangian describing this decay is [24]

$$\mathcal{L}_{\chi gg} \approx \frac{F\alpha_s}{4\sqrt{8}\,\pi}\,\frac{\sqrt{2\beta}}{m_\chi}\,\chi\, G^a_{\mu\nu} G^{a\,\mu\nu}\,, \tag{3.9}$$

where α_s is the strong gauge coupling constant and $F = \sum_f F_f$, see (3.5), with sum only over quarks. Higher order QCD corrections are also important here, and they coincide with those in the case of the SM Higgs boson. The produced gluons hadronise later on.

Obviously, well above the heavy quark threshold and well above QCD scale the description of the inflaton decays in terms of quarks and gluons is well-justified, while in the opposite case the effective description in terms of mesons is applicable. Between these ranges, where $m_\chi \simeq 1.5$–2.5 GeV, both approximations are not quite correct and no reliable description can be presented. At the same time, comparing relevant hadronic contributions described within these two approaches valid at lower and upper limits of this “untractable” interval, we observe deviations not larger than by an order of magnitude. Thus we conclude, that order-of-magnitude estimates of lifetime and decay rates of leptonic, photonic and total hadronic modes for the inflaton mass interval 1.5–2.5 GeV can be obtained by some interpolation between these “low-mass” and “high-mass” results.

Hence, for the inflaton decay rates to quarks we obtain the same formula as (3.6) multiplied by the number of colours, 3, and by a factor due to QCD corrections (having in mind uncertainties in the value of β discussed in section 2, for our estimates we adopt only

[4]We do not discuss decays below thresholds with one hadron being off-shell.

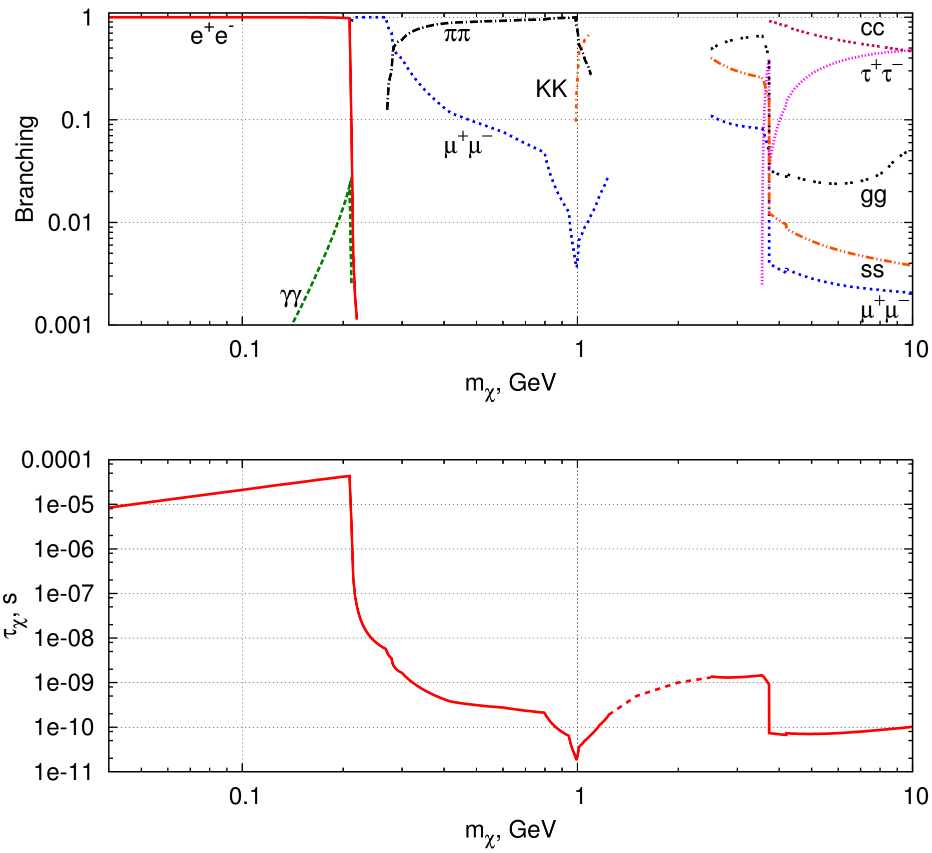

Figure 1. *Upper panel:* Inflaton branching ratios to various two-body final states as functions of inflaton mass m_χ, for 1.5-2.5 GeV range see discussion in the main text; *Lower panel:* inflaton lifetime τ_χ as a function of the inflaton mass m_χ. Here we set $\beta = \beta_0$, while for other values the life-time is inversely proportional to β. With β within the favourable interval (2.7) the lifetime can be two times smaller.

leading order QCD corrections presented in [25]). For decay rates into gluons we obtain

$$\Gamma_{\chi \to gg} = |F|^2 \left(\frac{\alpha_s}{4\pi}\right)^2 \frac{\beta\, m_\chi}{4\pi} ,$$

multiplied by the corresponding factor due to leading-order QCD-corrections [25].

The inflaton branching ratios and lifetime are presented in figure 1 as functions of the inflaton mass. Note, that inflaton partial widths into fermions decrease with increasing inflaton mass, while the opposite behaviour is observed for decays into pions due to the contribution of kinetic term from the trace of the energy-momentum. Diphoton mode varies with the inflaton mass and for the interesting range of parameters reaches its maximum of about 2.5% just at muon threshold. Thus, the sub-GeV inflaton predominantly decays into electrons, if its mass is below 200 MeV, otherwise to muons and pions with comparable branching ratios. If the inflaton mass is below the muon threshold, its life-time is about 10^{-5} s and above this threshold it falls rapidly down to 10^{-9} s. For the inflaton with mass well above 1 GeV, heaviest kinematically allowed fermions dominate the inflaton decay.

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Decays into photons and electrons are strongly suppressed for $m_\chi \gtrsim 1\,\mathrm{GeV}$, with branchings below 10^{-4}.

4 Inflaton from hadron decays

In this section we consider the inflaton production in rare decays of mesons and obtain corresponding limits on the light inflaton mass.

First, light inflaton can be produced in two-body meson decays. These are exactly the processes widely discussed in the past, when the SM Higgs boson was considered a (sub)GeV particle. Making use of the results from [14] one obtains for the amplitudes of the kaon decays

$$A\left(K^+ \to \pi^+\chi\right) \simeq \theta\frac{M_K^2}{v}\left\{\gamma_1\frac{1-\kappa}{2}\left(1-\frac{m_\chi^2-M_\pi^2}{M_K^2}\right)\right. \left. -\gamma_2(1-\kappa)+\frac{1}{2}\frac{3G_F\sqrt{2}}{16\pi^2}\sum_{i=c,t}V_{\mathrm{id}}^* m_i^2 V_{\mathrm{is}}\right\}, \tag{4.1}$$

$$A\left(K_L \to \pi^0\chi\right) = -A\left(K^+ \to \pi^+\chi\right)\ , \tag{4.2}$$

$$A\left(K_S \to \chi\pi^0\right) \simeq \theta\frac{3G_F\sqrt{2}}{16\pi^2}\frac{iM_K^2}{2v}\,\mathrm{Im}\sum_{i=c,t}V_{id}^* m_i^2 V_{is}\ , \tag{4.3}$$

with $\gamma_1 \sim 3.1\times 10^{-7}$ and much smaller γ_2. The largest contribution comes from the third term in (4.1). At quark level this term is due to the inflaton emission by a virtual quark in the quark-W-boson loop. For the branching ratios we get

$$\begin{aligned}\mathrm{Br}\left(K^+ \to \pi^+\chi\right) &= \frac{1}{\Gamma_{\mathrm{total}}(K^+)}\frac{|A(K^+\to\pi^+\chi)|^2}{16\pi M_K}\frac{2|\mathbf{p}_\chi|}{M_K}\\ &\approx 1.3\times 10^{-3}\cdot\left(\frac{2|\mathbf{p}_\chi|}{M_K}\right)\theta^2 \\ &\approx 2.3\times 10^{-9}\cdot\left(\frac{2|\mathbf{p}_\chi|}{M_K}\right)\cdot\left(\frac{\beta}{\beta_0}\right)\cdot\left(\frac{100\ \mathrm{MeV}}{m_\chi}\right)^2\ ,\end{aligned} \tag{4.4}$$

$$\begin{aligned}\mathrm{Br}\left(K_L \to \pi^0\chi\right) &\approx 5.5\times 10^{-3}\cdot\left(\frac{2|\mathbf{p}_\chi|}{M_K}\right)\theta^2\\ &\approx 1.0\times 10^{-8}\cdot\left(\frac{2|\mathbf{p}_\chi|}{M_K}\right)\cdot\left(\frac{\beta}{\beta_0}\right)\cdot\left(\frac{100\ \mathrm{MeV}}{m_\chi}\right)^2\ ,\end{aligned} \tag{4.5}$$

where $\mathbf{p}_\chi$ is the inflaton 3-momentum. The branching ratio of K_S is much smaller. For the inflaton mass in the kinematically allowed range the squared mixing angle θ^2 is of the order 10^{-5}–10^{-7} (2.4). For a model with $\beta = \beta_0$ branching ratios of the kaon decays are presented in figure 2 together with the relevant existing limits from the searches of the processes $K \to \pi + \mathrm{nothing}$ [26, 27]. It follows from figure 2 that the models with light inflatons,

$$m_\chi \lesssim 120\ \mathrm{MeV}\ ,$$

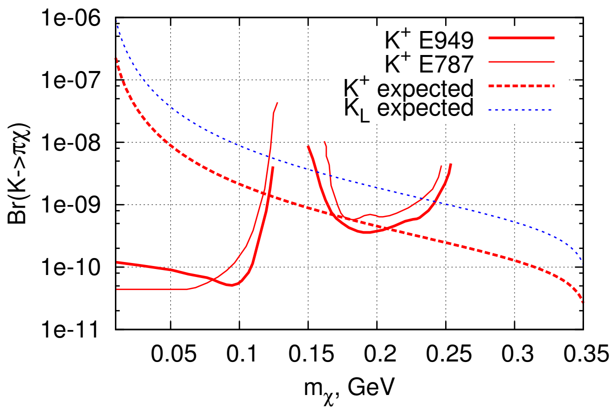

Figure 2. Expected branching ratios $\mathrm{Br}\,(K^+ \to \pi^+\chi)$, $\mathrm{Br}\,\left(K_L \to \pi^0\chi\right)$ and experimental bounds on $\mathrm{Br}\,(K^+ \to \pi^+\chi)$ from [27] and [26].

are *excluded* by negative results of these searches and models with 170 MeV $\lesssim m_\chi \lesssim$ 205 MeV are disfavoured. Further increase in sensitivity of these searches by one order of magnitude would allow to explore the light inflaton in the mass region 150–250 MeV. Thus, kaon decays are the most promising processes to search for the light inflaton.

In case of a larger inflaton mass heavy mesons have to be considered. The most promising here is the η-meson with the branching ratio of the order (cf. [14, 28])

$$\mathrm{Br}\left(\eta \to \chi\pi^0\right) \sim \frac{|\mathbf{p}_\chi|}{M_\eta} \cdot 10^{-6} \cdot \theta^2 \approx 1.8 \times 10^{-12} \cdot \left(\frac{2|\mathbf{p}_\chi|}{M_\eta}\right) \cdot \left(\frac{\beta}{\beta_0}\right) \cdot \left(\frac{100\ \mathrm{MeV}}{m_\chi}\right)^2 , \quad (4.6)$$

while the branching ratios of vector mesons are about two orders of magnitude smaller (cf. [29]).

Two-body decays of charmed mesons similar to $K \to \pi\chi$ are strongly suppressed as compared to that decay because of the smallness of the up-quark mass and the corresponding CKM matrix elements in the amplitude (cf. the third term in (4.1) which dominates). Three-body semileptonic decays have larger rates, however they are still quite small [30, 31],

$$\begin{aligned}
\mathrm{Br}\,(D \to e\nu\chi) &= \frac{\sqrt{2}G_F m_D^4}{96\pi^2 m_\mu^2 (1 - m_\mu^2/m_D^2)^2} \cdot \frac{7}{9}\theta^2 \cdot \mathrm{Br}(D \to \mu\nu) f\left(\frac{m_\chi^2}{m_D^2}\right) \\
&\approx 5.7 \times 10^{-9}\theta^2 f\left(\frac{m_\chi^2}{m_D^2}\right) , \\
f(x) &= (1 - 8x + x^2)(1 - x^2) - 12x^2 \log(x) .
\end{aligned}$$

Similar formula for the $D \to \mu\nu\chi$ decay rate is a bit more complicated because of the larger muon mass [30], but the rate is suppressed equally strong. On the contrary, decays of the beauty mesons are enhanced as compared to $K \to \pi\chi$. From ref. [32, 33] and (2.4),(2.6)

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we obtain for the light inflaton:

$$\begin{aligned}\mathrm{Br}\,(B \to \chi X_s) &\simeq 0.3 \frac{|V_{ts}V_{tb}^*|^2}{|V_{cb}|^2} \left(\frac{m_t}{M_W}\right)^4 \left(1 - \frac{m_\chi^2}{m_b^2}\right)^2 \theta^2 \\ &\simeq 10^{-6} \cdot \left(1 - \frac{m_\chi^2}{m_b^2}\right)^2 \left(\frac{\beta}{\beta_0}\right) \left(\frac{300\text{ MeV}}{m_\chi}\right)^2 ,\end{aligned} \tag{4.7}$$

where X_s stands for strange meson channel mostly saturated by a sum of pseudoscalar and vector kaons. The inflaton with the mass below the muon threshold escapes the detector (see figure 1) giving the signatures

$$B \to K + \text{nothing} , \qquad B \to K^* + \text{nothing} .$$

Heavier inflaton can decay within the detector, with the most clear mode being the muon one at the level of 0.01–1, depending on the mass, see figure 1. Note, that collected world statistics at b-factories allowed to measure branching fractions of relevant decays $B \to K^{(*)} l^+ l^-$ with accuracy of about 10^{-7} [34], which is comparable to the expected signal (4.7). Thus, an appropriate reanalysis of these data might give a chance to probe the inflaton of mass $m_\chi \sim 300\,\text{MeV}$.

Inflaton can also be produced in other meson decays and heavy baryon decays. We do not discuss these channels considering them as subdominant for the inflaton production and less promising in direct searches for the light inflaton.

5 Inflaton production in particle collisions

In this section we discuss in detail the inflaton production in particle collisions. If the collision energy is large enough, the most efficient mechanism of inflaton production is kinematically allowed decays of heavy mesons, produced in the collision. Then the production cross section σ can be estimated as the product of the meson production cross section and the branching of the meson decay into the inflaton,

$$\frac{\sigma}{\sigma_{pp,\text{total}}} = M_{\text{pp}} \big(\chi_s (0.5\,\mathrm{Br}(K^+ \to \pi^+ \chi) + 0.25\,\mathrm{Br}(K_L \to \pi^0 \chi)) + \chi_c\,\mathrm{Br}(B \to \chi X_s)\big) , \tag{5.1}$$

where total hadron multiplicity M_{pp} and relative parts going into different flavours $\chi_{s,c}$ are given in table 1 for several existing beams, and $\sigma_{pp,\text{total}} \simeq 40$ mbarn is the total proton cross section [10]. Here only strange and beauty mesons were taken into account, as they give the main contribution. In figure 3 we present the estimate of this indirect inflaton production in a beam-target experiment for several available proton beams.

In the models with the inflaton mass above the bottom quark threshold, the dominant source of the inflatons is the direct production in hard processes similar to the case of the SM Higgs boson. The main channel for the inflaton is the same as for the Higgs boson, i.e. the gluon-gluon fusion [38]. The inflaton production is calculated exactly as the production of the Higgs boson of the same mass, only it is suppressed by the mixing angle squared θ^2. For several available proton beams the result of the calculation with the

Experiment	E, GeV	$N_{\rm POT}$, 10^{19}	$M_{\rm pp}$ [10]	χ_s [35, 36]	χ_c [37]	χ_b [37]
CNGS [4]	400	4.5	13	1/7	$0.45 \cdot 10^{-3}$	$3 \cdot 10^{-8}$
NuMi [3]	120	5	11	1/7	$1 \cdot 10^{-4}$	10^{-10}
T2K [2]	50	100	7	1/7	$1 \cdot 10^{-5}$	10^{-12}
NuTeV [5]	800	1	15	1/7	$1 \cdot 10^{-3}$	$2 \cdot 10^{-7}$

Table 1. Adopted values of relevant for inflaton production parameters of several experiments.

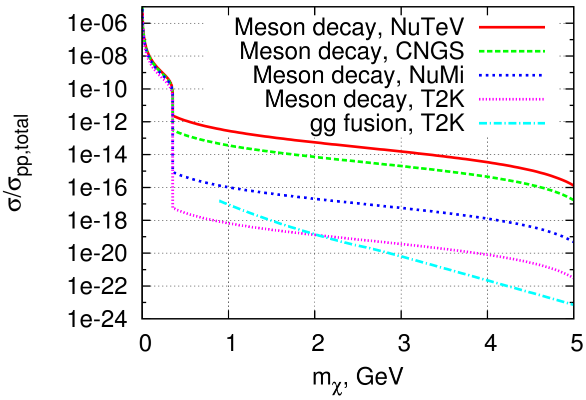

Figure 3. The inflaton production cross section in pp collisions normalised to the total pp cross section $\sigma_{pp\ total} \simeq 40$ mbarn. For all graphs $\beta = 1.5 \times 10^{-13}$. The lines are for several existing beam energies (see table 1). For the T2K case the estimate of the direct production by gluon fusion is also given, while for higher beam energies it is negligible, compared to the meson decay channel of production. At the kaon threshold $m_\chi \sim m_K$ the account of the contribution from the η-meson decay (4.6) might somehow smooth the cross section change.

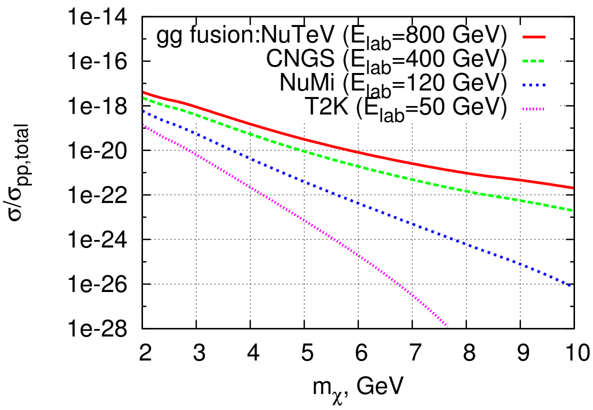

Figure 4. The ratio of the inflaton direct production cross section in pp collisions to the total pp cross section, $\sigma_{pp\ total} \sim 40$ mbarn, for inflationary $\beta = \beta_0$.

parton distribution functions from [39] is given in figure 4. Here we present the estimate of the direct contribution to the inflaton production even for small inflaton mass down to $m_\chi \gtrsim 2\,$GeV, to illustrate the statement that indirect production dominates below the B-meson threshold, cf. figure 3.

6 Limits from direct searches and predictions for forthcoming experiments

The light inflaton will be produced in large amounts in any beam dump, and then decay further into a photon or lepton pair, depending on its mass. Search of a penetrating particle of this type was performed by the CHARM experiment [40]. In this article the search for a generic axion was performed. However, the only difference between the axion in [40] and the light inflaton is in the estimate of the decay rate and production cross section. Here we will make rough reanalysis for the case of the light inflaton. As far as in the most interesting region of masses the inflaton lifetime is small and a significant amount of inflatons decay before they reach the detector, this reanalysis can not be made by simple rescaling of the result. So, first we will reproduce in the simplest possible way the resulting bound from [40], and then change in the analysis the decay and production rates. The resulting picture in [40] can be reproduced by demanding that the number of decays in the detector is larger than the background. The number of decays in the detector is roughly estimated as

$$N \simeq N_0 \sigma_X \mathrm{e}^{-\Gamma \frac{l_{\mathrm{dec}}}{\gamma}} \left(1 - \mathrm{e}^{-\Gamma \frac{l_{\mathrm{detector}}}{\gamma}}\right) , \tag{6.1}$$

where N_0 is the overall coefficient describing luminosity, σ_X is the production cross section of the axion (eq. (5) in [40]), Γ is the decay width (sum of eqs. (3) and (4) for muons and electrons in [40]), $l_{\mathrm{dec}} = 480$ m is the decay length before the detector, $l_{\mathrm{detector}} = 35$ m is the detector length, $\gamma = E/m_X$ is the typical relativistic gamma factor of the axion, with $E \sim 10$ GeV. Then, figure 4 from [40] is approximately reproduced for $N/(N_0\sigma_{\pi^0}) \simeq 10^{-17}$.

Using the same logic we can obtain the bound for the inflaton. We then get instead of (6.1)

$$N \simeq N_0 \sigma_{\mathrm{prod}} (\mathrm{Br}(\chi \to \gamma\gamma) + \mathrm{Br}(\chi \to ee) + \mathrm{Br}(\chi \to \mu\mu)) \\ \times \mathrm{e}^{-\Gamma \frac{l_{\mathrm{dec}}}{c\gamma}} \left(1 - \mathrm{e}^{-\Gamma \frac{l_{\mathrm{detector}}}{c\gamma}}\right) , \tag{6.2}$$

where for the inflaton production cross section σ_χ we take (5.1), and we adopt the simple estimate of the π_0 yield: $\sigma_{\pi^0}/\sigma_{pp,\mathrm{total}} = M_{\mathrm{pp}}/3$, with $M_{\mathrm{pp}} = 13$ for the CNGS beam (see table 1). Then, the region forbidden by the experiment [40] is given in figure 5.

One can see, that for the inflationary self coupling $\beta = \beta_0$ the CHARM results exclude the inflaton with masses $m_\chi \gtrsim 280\,\mathrm{MeV}$, while for the upper limit of the interval (2.7) we have

$$m_\chi \gtrsim 270 \text{ MeV} . \tag{6.3}$$

One should be somewhat careful taking this limit. The exact dependence of the bound on β in this region is sensitive to the proper analysis of the CHARM experiment (proper simulation of energy distribution of the produced inflatons, detector sensitivity to different decay modes, etc.). From the plot in figure 5 we conclude, however, that the value of $m_\chi > 210$ MeV is a conservative bound: within the simple approach above we get an order of magnitude "safety margin". The definite conclusion about higher values of masses for $\beta \sim \beta_0$ requires careful reanalysis of the CHARM data.

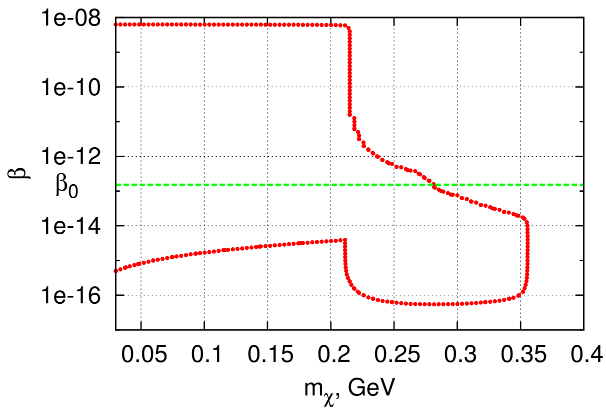

Figure 5. Limits on the inflaton properties from the CHARM axion searches. The region of (m_χ, β) surrounded by the curve is forbidden. One can see, that for inflationary $\beta = \beta_0$ the inflatons with masses up to approximately 280 MeV are ruled out.

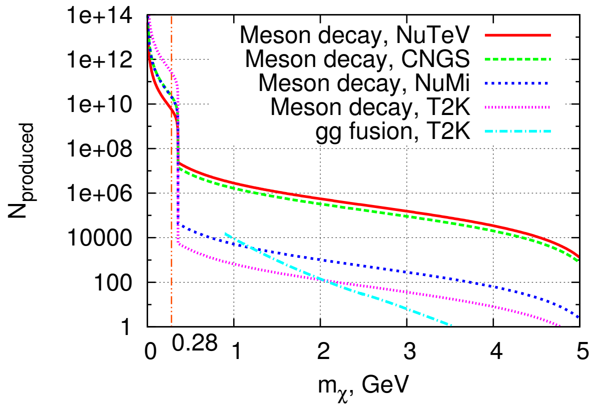

Figure 6. Number of the inflatons produced with one-year statistics for the reference beams. All geometric factors are neglected. The vertical line indicates the lower limit on the inflaton mass protect(6.3).

For the allowed inflaton mass range and using the results obtained in section 5 we estimate the number of inflatons produced during one year of running at designed luminosity for several experiments by multiplying the cross section ratio by the number of protons on target, see figure 6 and table 1. This number should be taken with a grain of salt, as far as we have totally neglected all possible geometrical factors and possible collimation or deflection of the produced charged kaons. However one can conclude that higher energy beams are certainly preferable in the searches for the inflaton, because the dominant production mode is decays of beauty hadrons.[5] For small mass the number of the inflatons per year can exceed several millions, while at $m_\chi \sim 5\,\text{GeV}$ it is about a thousand at best.

[5]This is not so for low mass inflatons, which can be produced in kaon decays and one can win slightly because of usually higher luminosity of the lower energy beam.

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7 Conclusions

In this paper we presented an example of a simple inflationary model, which can be fully explored in particle physics experiments. The inflaton is light and can be hunted for in the beam-target experiments and in two-body meson decays. The inflaton life-time is about 10^{-9}–10^{-10} s, so it decays close to the production point. The dominant mechanism of inflaton production is the meson decays. Thus, and it is quite amusing in fact, the first stage of the Universe evolution described by means of quantum field theory can be directly tested with modern experimental techniques.

The mass limits on the light inflaton are found to be $270\text{ MeV} \lesssim m_\chi \lesssim 1.8\,\text{GeV}$. Note, that careful reanalysis of the results of the CHARM experiment [40] may change the lower bound given here.

The paper analysed the quartic inflaton self-interaction with minimal (or very weakly non-minimal) coupling to gravity, which is still allowed by the current WMAP data. In near future results from the PLANCK experiment will determine the inflationary parameters — spectral index n_s and tensor-to-scalar ratio r—with much better precision. This will allow to fix the value of the non-minimal coupling constant ξ (see relations between ξ and n_s, r in [18–20]). In case it is large, the laboratory detection will still be possible, while the analysis of the current work should be rescaled for larger value of the self-coupling constant β and thus for stronger signal.

Note also that the inflaton model discussed in this paper has been suggested in ref. [16] as an extension of the νMSM [41, 42]. In that joint model the inflaton vacuum expectation value provides the three sterile neutrino of νMSM with Majorana masses, while decays of inflatons in the early Universe to the lightest sterile neutrino produce the dark matter. In appendix A we check, that the limits presented in νMSM itself do not limit the interval for the inflaton mass obtained in this paper. On the contrary, the upper limit on the inflaton mass constrains the mass of dark matter sterile neutrino in this model. This proves that inflation can be directly tested in fully realistic extensions of the SM.

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A The νMSM extension

In ref. [16] the inflaton model considered in this paper has been implemented in the framework of νMSM [41, 42].

Neutrino oscillations is the only direct evidence of physics beyond the Standard Model of particle physics. Cosmology provides two other evidences — dark matter and baryon asymmetry of the Universe — which are not explained within SM if General Relativity is a correct theory of gravity. These three problems can be addressed within a neutrino minimal Standard Model (νMSM) [41, 42], suggested as a minimal version of SM capable of explaining all of the three. Enlarged by the additional scalar field [16] coupled to the SM Higgs boson via tree-level-scale-invariant interaction, νMSM provides both early-time inflation and common source of electroweak symmetry breaking and of sterile neutrino masses. This one-energy scale model[6] is a minimal, full and self-consistent extension of the Standard Model of particle physics.

This extension implies addition to the model (2.1) three new singlet right-handed neutrinos N_I, $I = 1, 2, 3$ with the Lagrangian

$$\mathcal{L}_{\nu\text{MSM}} = i\bar{N}_I \gamma_\mu \partial^\mu N_I + \left(F_{\alpha I} \bar{L}_\alpha N_I \tilde{H} - \frac{f_I}{2} \bar{N}_I^c N_I X + \text{h.c.} \right) ,$$

where L_α ($\alpha = 1, 2, 3$) are charged lepton doublets and $\tilde{H} = \epsilon H^*$, where ϵ is 2×2 antisymmetric matrix and H is the Higgs doublet.

To begin with, let us discuss the allowed ranges of the Yukawa coupling constants f_I (Yukawas $F_{I\alpha}$ are irrelevant for this study). The flatness of the inflaton potential implies smallness of the quantum corrections to the quartic coupling (2.6). Requiring again smaller than 10% contribution we obtain

$$f_I < 1.5 \times 10^{-3} \tag{A.1}$$

and get for the sterile neutrino masses

$$M_I = f_I \langle X \rangle < 270 \cdot \left(\frac{m_\chi}{100 \text{ MeV}} \right) \left(\frac{1.5 \times 10^{-13}}{\beta} \right)^{1/2} \left(\frac{f_I}{1.5 \times 10^{-3}} \right) \text{ GeV} . \tag{A.2}$$

A lower limit on the masses of the two heavier sterile neutrinos, $M_{2,3}$, follows from successfulness of the Big Bang Nucleosynthesis (BBN), that does not constrain the inflaton mass.

The lightest sterile neutrino in the model (2.1) can be stable at cosmological time-scale and comprise the dark matter of the Universe. In this case, the sterile neutrino is quite light, $M_1 \lesssim 1\,\text{GeV}$, and should not get thermalised in primordial plasma, to be viable dark matter. The latter is natural, since stability at cosmological time-scale implies extremely small (if any) mixing with active neutrino via $F_{1\alpha}$, and small mass implies tiny coupling to inflaton. Light inflatons are in equilibrium in the primordial plasma and decay to the lightest sterile neutrino mostly at the temperature $T \sim m_\chi$. They provide contribution of the sterile neutrino N_1 to the energy density of the Universe today [16]

$$\Omega_N = \frac{1.6 f(m_\chi)}{S} \cdot \frac{\beta}{1.5 \times 10^{-13}} \cdot \left(\frac{M_1}{10 \text{ keV}} \right)^3 \cdot \left(\frac{100 \text{ MeV}}{m_\chi} \right)^3 , \tag{A.3}$$

where $S > 1$ is a dilution factor accounting for a possible entropy production due to the late decays of the heavier sterile neutrinos [44] and the function $f(m_\chi)$ is determined by the number of degrees of freedom $g_*(T)$ in a primordial plasma at inflaton decays. It changes monotonically from 0.9 to 0.4 for inflaton mass from 70 MeV to 500 MeV and for

[6] For a discussion of gauge hierarchy problem in models of this type see, e.g. ref. [43].

heavier inflaton can be approximated as $f\left(m_\chi\right) \simeq \left[10.75/g_*\left(m_\chi/3\right)\right]^{3/2}$. Other mechanisms can also contribute to the dark matter production. Hence, at a given m_χ equation (A.3) implies the *upper limit*

$$\frac{M_1}{10\ \text{keV}} < \left(\frac{S}{6.4\, f(m_\chi)}\right)^{1/3} \left(\frac{\Omega_N}{0.25}\right)^{1/3} \left(\frac{1.5\times10^{-13}}{\beta}\right)^{1/3} \left(\frac{m_\chi}{100\ \text{MeV}}\right). \tag{A.4}$$

The corresponding upper limit on f_1 supersedes (A.1).

There are several notes in order. First, from (A.4) one concludes, that with the inflaton mass in the range (2.8), (2.13), the lightest sterile neutrino can *naturally* be a Warm Dark Matter candidate. Second, with model parameters tuned within their ranges to minimise r.h.s. of the inequality (A.4), one does not exceed the lower bound on the sterile neutrino mass, $M_1 \gtrsim 1.7$ keV, from the study [45, 46] of the dark matter phase space density in dwarf spheroidal galaxies. Hence, in this model limits from the dark matter sector do not shrink the allowed parameter region, and in particular, the inflaton mass range. Third, maximising r.h.s. of inequality (A.4) one obtains an *upper limit* on the lightest sterile neutrino mass,

$$M_1 \lesssim 13 \cdot \left(\frac{m_\chi}{300\ \text{MeV}}\right) \left(\frac{S}{4}\right)^{1/3} \cdot \left(\frac{0.9}{f\left(m_\chi\right)}\right)^{1/3} \text{keV}\,. \tag{A.5}$$

Hence, given the range (2.15) the lightest sterile neutrino is significantly lighter than electron (and other charged SM fermions). Since the light inflaton decays to SM fermions due to mixing with the Higgs boson, these partial decay rates are proportional to the squared masses of the corresponding fermions similar to its decay rates to sterile neutrinos. The lightness of the dark matter sterile neutrino guarantees, that inflaton decays to the dark matter neutrino never suppresses its decay branching ratios to visible channels.

From the formulas above one concludes that the limits on the inflaton mass obtained in this paper are not affected due to additional constraints typical for νMSM. At the same time, the limits (A.5), (2.15) imply that the dark matter is lighter than about 100 keV and with account of the allowed Higgs boson mass (see section 2). Once the inflaton is found the upper limit (A.5) on the dark matter neutrino mass will be settled. And vice versa, once the dark matter neutrinos are found, eq. (A.5) fixes a *lower* limit on the inflaton mass; this limit can supersede limit (6.3) from direct searches, if $M_1 \gtrsim 10$ keV.

Note that inflaton can decay into sterile neutrinos $N_{2,3}$ if kinematically allowed. The decay mode into the not-dark-matter sterile neutrinos $N_{2,3}$ (invisible mode) has the width

$$\Gamma_{\chi\to N_I N_I} = \frac{\beta M_I^2}{8\pi m_\chi}\left(1-\frac{4m_f^2}{m_\chi^2}\right)^{3/2}.$$

This formula is very similar to (3.6). This decay mode can even be dominant, if the sterile neutrinos $N_{2,3}$ are the heaviest fermions between the kinematically allowed ones. As far as $M_{2,3} \lesssim m_\pi$ is disfavoured by the BBN, this can be relevant for inflaton masses above approximately 300 MeV. As a consequence, inflaton lifetime can be somewhat shortened as compared to the results presented in figure 1b.

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