

Inflation might be caused by the right

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Inflation might be caused by the right (handed neutrino)

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ABSTRACT: We show that the scalar field that drives inflation can have a dynamical origin, being a strongly coupled right handed neutrino condensate. The resulting model is phenomenologically tightly constrained, and can be experimentally (dis)probed in the near future. The mass of the right handed neutrino obtained this way (a crucial ingredient to obtain the right light neutrino spectrum within the see-saw mechanism in a complete three generation framework) is related to that of the inflaton and both completely determine the inflation features that can be tested by current and planned experiments.

KEYWORDS: Cosmology of Theories beyond the SM, Beyond Standard Model, Neutrino Physics

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1 Introduction

Modern cosmology is built upon a theoretical framework whose foundations are Big Bang theory and general relativity. Together they describe with precision many aspects of the universe from the first nanosecond until our days. Despite its power, this framework is unable to explain the observed flatness and homogeneity of space, let alone the origin of matter and structure. As a consequence, this minimal picture is usually supplemented by postulating one or more episodes of accelerated expansion [1] (inflation).

The simplest realization of this idea is an approximately constant energy density, leading to quasi-de-Sitter or exponential expansion. This can be parametrized in terms of a fundamental scalar field, named the inflaton, whose nature is yet unknown. Any other species that might have been present together with the inflaton are quickly diluted away by the expansion, so that the inflaton is essentially playing “solo” during the inflationary epoch.

Although this idea is indeed attractive, fundamental scalars have not been observed yet. Even more, since the rise and fall of the aether — a primitive version of a fundamental scalar — alternative scenarios have been explored, in which the scalar is an order parameter of some strong dynamics, rather than a fundamental degree of freedom. In the famous BCS or Nambu-Jona-Lasinio mechanisms [2] new interactions associated with a high energy scale Λ are used to trigger the formation of a low energy condensate, which mimics the role of a scalar.

This idea was followed by Bardeen, Hill and Lindner [3], who proposed that a top quark condensate can replace the fundamental standard model Higgs boson to drive electroweak symmetry breaking. A four fermion self-coupling of the top quark of strength G signals the formation of a top-antitop condensate, dynamically generating a mass for the top. Below the cutoff scale Λ , one can integrate out the high frequency modes of the fermions, obtaining an effective theory of a Higgs-like composite. In the large N_c limit the theory predicts both the top mass and the scale of electroweak symmetry breaking in terms of the

fundamental parameters of the fermion theory at the cutoff scale. The scale of electroweak symmetry breaking can be parametrically small compared to Λ .

In this work, we will play the same game with the inflaton. We will attribute the dynamical origin of the inflaton field, to another “solo” player, the right handed neutrino [4]. In analogy to the idea of top-quark condensation, we consider the standard model including right-handed neutrinos but without an inflaton. We add four-fermion couplings which should be viewed as effective interactions that describe the physics below a high energy cutoff, that we will take to be the Planck scale. This new interaction should be strong enough for a neutrino condensate that will trigger spontaneous symmetry breaking of lepton number and produce a Majorana mass for the right-handed neutrino. The same dynamics also produces “natural inflation” [5]. At the same time that we obtain a slightly red spectral index, both the inflationary de Sitter scale and the right-handed neutrino mass can naturally be of order 10^{17} GeV. As compared to the usual approach to inflaton model building our dynamical framework is both economical and predictive.

2 Constructing the scalar field

For simplicity we assume that one generation of right handed neutrinos has a four-fermion self-coupling, generated at a high energy scale Λ from unknown dynamics. The underlying new dynamics might for example be some non-abelian gauge interactions. The four-fermion effective interaction for the right handed neutrinos below the scale Λ takes the form

$$G (\bar{\nu}_R^c \nu_R) (\bar{\nu}_R \nu_R^c) \tag{2.1}$$

where G is the dimensionful coupling constant, ν_R is the right handed neutrino and ν^c indicates charge conjugation. This is an effective interaction describing the physics below the physical cutoff Λ . There may be other higher dimension operators, but these will have subdominant effects at energies substantially below the cutoff scale.

Analogously to [3] for a top condensate, the gap equation for a dynamically generated right-handed neutrino mass has a solution when

$$G\Lambda^2 > \frac{8\pi^2}{N_f}, \tag{2.2}$$

where N_f is the number of right-handed neutrino flavors. If the four-fermion interaction were the result of integrating out a new non-abelian gauge interaction, N_f could be the number of “colors” of this gauge theory. Strictly speaking, the analysis of the condensate properties is performed in the limit of large N_f .

When the right-handed neutrinos condense the condensation effects can be incorporated by introducing an auxiliary scalar field Φ into the lagrangian

$$- m_o^2 \Phi^\dagger \Phi + g_o (\bar{\nu}_R^c \nu_R \Phi + \text{h.c.}). \tag{2.3}$$

Notice that the new effective scalar field does not have a kinetic term and reproduces the four-fermion interaction when integrated out with

$$G = g_o^2 / m_o^2. \tag{2.4}$$

For the study of the low-energy regime we would like to keep the effective scalar field and integrate out the short distance components of the right handed neutrino field. By doing so, we will see that at scales below the cutoff, the effective scalar field develops fully gauge invariant induced kinetic terms and quartic self-interactions through fermion loops. The full induced effective lagrangian will take the form

$$g_o (\bar{\nu}_R^c \nu_R \Phi + \text{h.c.}) + Z_\Phi |D_\mu \Phi|^2 - m_\Phi^2 \Phi^\dagger \Phi - \lambda_o (\Phi^\dagger \Phi)^2, \quad (2.5)$$

where D_μ is the gauge covariant derivative, and all loops now to be defined with respect to a low energy scale μ yielding for the induced parameters

$$Z_\Phi = \frac{N_f g_o^2}{(4\pi)^2} \ln \left(\frac{\Lambda^2}{\mu^2} \right) \quad (2.6)$$

$$m_\Phi^2 = m_o^2 - \frac{2 N_f g_o^2}{(4\pi)^2} (\Lambda^2 - \mu^2) \quad (2.7)$$

$$\lambda_o = \frac{N_f g_o^4}{(4\pi)^2} \ln \left(\frac{\Lambda^2}{\mu^2} \right). \quad (2.8)$$

The mechanism for spontaneous symmetry breaking is now apparent in the effective scalar mass-squared, which is shifted by a negative finite value proportional to the cutoff Λ .

We would like to get a Lagrangian with a canonical kinetic term, which can be done by rescaling the scalar field $\Phi \rightarrow \Phi/\sqrt{Z_\Phi}$ and defining

$$g = g_o/\sqrt{Z_\Phi} \quad (2.9)$$

$$m^2 = m_\Phi^2/Z_\Phi \quad (2.10)$$

$$\lambda = \lambda_o/Z_\Phi^2, \quad (2.11)$$

to get

$$g (\bar{\nu}_R^c \nu_R \Phi + \text{h.c.}) + |D_\mu \Phi|^2 - V(\Phi), \quad (2.12)$$

where $V(\Phi)$ is the scalar field potential and is given by

$$V(\Phi) = m^2 \Phi^\dagger \Phi + \lambda (\Phi^\dagger \Phi)^2. \quad (2.13)$$

For G satisfying (2.2), m^2 becomes negative at sufficiently low energies, leading to spontaneous symmetry breaking with Φ developing a vev $v = \sqrt{-m^2/\lambda}$.

In the cosmological context we will want the analog of the effective potential (2.13) at finite temperature; this is obtained by replacing (2.7) by

$$m_\Phi^2 = m_o^2 + k^2 T^2 - \frac{2 N_f g_o^2}{(4\pi)^2} (\Lambda^2 - k^2 T^2), \quad (2.14)$$

and similarly replacing μ by kT in (2.6) and (2.8). When G satisfies (2.2), we can simplify the notation by defining

$$\delta \equiv \frac{N_f G \Lambda^2}{8\pi^2} - 1, \quad (2.15)$$

$$\beta \equiv \frac{m_o}{\Lambda}. \quad (2.16)$$

At finite temperature the effective potential acquires a symmetry-breaking minimum below the critical temperature given by

$$k^2 T_c^2 = \frac{\delta\beta^2 \Lambda^2}{1 + (1 + \delta)\beta^2}. \quad (2.17)$$

The vacuum expectation value of Φ , which we denote by v , can be written as

$$v^2 = \frac{\delta\beta^2 \Lambda^2}{\lambda_0} \left(1 - \frac{T^2}{T_c^2}\right) \quad (2.18)$$

$$= \frac{g^2 \delta\beta^2}{g_0^4} \Lambda^2 \left(1 - \frac{T^2}{T_c^2}\right). \quad (2.19)$$

3 The phase field and its circumstances

In the previous section we have generated a potential for the effective scalar field that represents the physics of the neutrino condensate. For cosmology we will introduce by hand a cosmological constant term so that $V(\Phi = \sqrt{-m^2/\lambda}) = 0$. With this addition, the potential can now be written as

$$V(\Phi) = \lambda \left(\Phi^\dagger \Phi - \tilde{m}^2\right)^2 \quad (3.1)$$

with $\lambda \tilde{m}^2 = m^2$ and $\lambda \tilde{m}^4$ the cosmological constant term mentioned before. Notice that this potential is invariant under a global U(1) transformation $\Phi \rightarrow e^{i\alpha} \Phi$, which is nothing less than lepton number.

Since Φ is a complex scalar field, it can be parametrized as $\Phi = \phi e^{i\theta}$, and the induced potential $V(\Phi)$ is a function of the radial field only, i.e.

$$V(\Phi) = V(\phi) = \lambda (\phi^2 - \tilde{m}^2)^2. \quad (3.2)$$

The radial field has a mass $m_\phi^2 = \lambda \tilde{m}^2$.

Potentials of the form (3.2) are easy to analyze regarding inflation: in order to obtain sufficient inflation (the famous 60 e-folds), the initial value of the field ϕ must be greater than the Planck mass. In order to obtain acceptable density perturbations, λ must be about 10^{-15} [6]. Such a small value for λ can be considered the kiss of death regarding the use of the radial field potential for inflation. In our model, λ is generated dynamically, not chosen by hand; from (2.11) it can be clearly seen that there is no $\Lambda - \mu$ combination for which λ can be so miserably small.

The phase field θ on the other hand, is a Nambu-Goldstone boson and is massless at tree level. If the U(1) symmetry of the potential is preserved, the phase field will remain massless even with loop corrections. If the U(1) is explicitly broken, the field θ will acquire a potential from loop corrections leading to nonzero mass, becoming a pseudo Nambu-Goldstone boson. If this new mass is hierarchically smaller than that of the radial field, θ will be effectively massless near the original symmetry breaking scale. When the temperature of the thermal bath drops below the mass of the radial field, its excitations

will be damped, and its expectation value rapidly settles into the temperature-dependent minimum given by (2.18). Thus for temperatures less than $O(T_c)$ the condensate dynamics is described by θ alone. Its potential $V_\theta(\theta)$, which was negligible for high temperatures, becomes important below T_c . This phase field then rolls down its potential to its minimum, acting as the driver for an extended period of inflation.

To compute the outcome of inflation driven by a neutrino condensate, we first need to specify the explicit breaking of the U(1) symmetry from dynamics above the cutoff scale. On general grounds it is expected that global symmetries such as our lepton number might be broken explicitly by Planck scale physics [7]. Two general arguments support this assertion. The first comes along the black-hole “no hair” theorems of classical relativity: as black-holes cannot support any global “hair”, a virtual exchange of black-holes would give rise to global symmetry violating operators in the low energy theory, where low energy in this case means energies smaller than the mass of the black-hole. The second argument arises courtesy of the existence of wormhole configurations [8]. While gauge charges cannot be sent down a wormhole there is nothing to prevent the loss of global charges this way [9]. The effective theory at scales below the wormhole scale will have non-zero global charge carrying operators. Integrating out the effective black holes will bring Planck scale suppressed non-renormalizable operators that break the global U(1).

Thus if we take the cutoff scale to be the Planck scale, $\Lambda \simeq 10^{19}$ GeV, it is natural to expect an explicit breaking of the global U(1). The lowest dimension symmetry-breaking operator constructed from the right-handed neutrinos is given by

$$G' \left[(\bar{\nu}_R^c \nu_R)^2 + (\bar{\nu}_R \nu_R^c)^2 \right]. \tag{3.3}$$

As we have seen before, we can reproduce the physics of the four-fermion coupling by resorting to the scalar field Φ as before, adding the interaction

$$g' \left(\bar{\nu}_R^c \nu_R \Phi^\dagger + \bar{\nu}_R \nu_R^c \Phi \right). \tag{3.4}$$

Under a global phase redefinition,

$$\begin{aligned} \nu_R &\longrightarrow e^{i\alpha} \nu_R \\ \nu_R^c &\longrightarrow e^{-i\alpha} \nu_R^c \\ \Phi &\longrightarrow e^{2i\alpha} \Phi \end{aligned}$$

the new term is not invariant, explicitly breaking the U(1) symmetry down to a residual discrete symmetry generated by $\theta \rightarrow \theta + 2\pi$. This results in a nontrivial potential for the effective pseudo Nambu-Goldstone boson from the condensate. This field will play the role of the inflaton in our model.

Neglecting the shift in the vacuum expectation value of the radial field induced at 1-loop, the tree level right handed neutrino mass will now be given by

$$m_R^2(\theta) = (g^2 + g'^2 + 2gg' \cos(\theta))v^2, \tag{3.5}$$

with v^2 given by (2.18), in the approximation that $g' \ll g$. The explicitly broken U(1) symmetry is reflected in the mass dependence of the right handed neutrino mass on the

phase field θ . Quantum effects will produce a potential for this field, thus providing our inflationary potential. At the 1-loop level this is given by

$$\begin{aligned}
 V_\theta(\theta) &= -\frac{1}{(16\pi)^2} (m_R^2(\theta))^2 \ln\left(\frac{m_R^2(\theta)}{v^2}\right) \\
 &= -\frac{g^2 g'^2 v^4}{16\pi^2} \left[\frac{g^2 + g'^2}{2gg'} + \cos(\theta)\right]^2 \ln\left[2gg' \left(\frac{g^2 + g'^2}{2gg'} + \cos(\theta)\right)\right]. \quad (3.6)
 \end{aligned}$$

This potential has extrema at $\theta = 0, \pi$, $\cos^{-1}\left(-\frac{g^2 + g'^2}{2gg'} + \frac{1}{2gg'\sqrt{e}}\right)$ being $\theta = \pi$ the only minimum.

The mass of the pseudo Nambu-Goldstone boson at the minimum of the potential, i.e. the mass of the inflaton is given by

$$m_\theta^2 \equiv \left. \frac{\partial^2 V_\theta}{\partial \theta^2} \right|_{\theta=\pi} = -\frac{gg'v^2}{32\pi^2} (g - g')^2 \left[1 + 2\ln\left((g - g')^2\right)\right] \quad (3.7)$$

while that of the right handed neutrino reads

$$m_R^2 \Big|_{\theta=\pi} = (g - g')^2 v^2 \quad (3.8)$$

Notice that the potential as well as both masses are invariant under the exchange $g \longleftrightarrow g'$ and that both masses vanish when $g = g'$; this corresponds to the degenerate case that only the real part of Φ couples to the neutrinos, meaning that the condensate is uncharged.

As in the radial field case, the potential for the phase field does not vanish at its minimum, a fact that we are going to change by defining a new phase field potential given by

$$\begin{aligned}
 V(\theta) &= V_\theta(\theta) - V_\theta(\theta = \pi) \\
 &= -\frac{g^2 g'^2 v^4}{16\pi^2} \left\{ \left[\frac{g^2 + g'^2}{2gg'} + \cos(\theta)\right]^2 \ln\left[2gg' \left(\frac{g^2 + g'^2}{2gg'} + \cos(\theta)\right)\right] \right. \\
 &\quad \left. - \left[\frac{(g^2 - g'^2)^2}{2gg'}\right]^2 \ln\left[(g - g')^2\right] \right\}. \quad (3.9)
 \end{aligned}$$

On general grounds, one would expect $g' \ll g$, i.e. small explicit breaking. In this case, the effective potential takes the delightfully simple form

$$V(\theta) \simeq -\frac{g^3 g' v^4}{32\pi^2} [1 + 2\ln(g^2)] (1 + \cos(\theta)) \quad (3.10)$$

which is of the form of the well-known natural inflation potential [5] $M^4 (1 + \cos(\theta))$ with

$$M^4 = -\frac{g^3 g' v^4}{32\pi^2} [1 + 2\ln(g^2)]. \quad (3.11)$$

Notice that although *prima facie* this constant term looks negative, it is indeed positive for $g < 1$.

4 Inflation phenomenology

Detailed analysis of the virtues of natural inflation models already exist in the literature [10]. Here we will briefly review the basic features of our model and derive the constraints that available data impose on the model parameters.

For an inflationary theory to correspond to the observed universe, it must satisfy at least two conditions: (i) it has to explain the observed thermal equilibrium of the cosmic microwave background radiation (CMB) which is guaranteed by providing sufficient inflation, i.e. the inflationary potential must drive an increase on the scale factor of a minimum of e^{60} ; (ii) the quantum fluctuations of the inflaton should give rise to primordial density fluctuations of size $\delta\rho/\rho$ and spectral index n_s in agreement with observations [11].

During inflation, the inflaton rolls down towards the minimum of its potential, evolving according to

$$\ddot{\theta} + 3H\dot{\theta} + \partial V/\partial\theta = 0 \tag{4.1}$$

where the Hubble rate H is given by

$$H^2 = \frac{8\pi}{3m_{\text{Pl}}^2} \left[\frac{\dot{\theta}^2}{2} + V(\theta) \right]. \tag{4.2}$$

In general, we are interested in potentials which contain one region flat enough that the evolution of the field is friction dominated, (what goes under the name of slow-roll approximation) so that the equation of motion is essentially given by

$$3H\dot{\theta} + \partial V/\partial\theta = 0. \tag{4.3}$$

Within the slow-roll approximation, the number of e-folds of inflation when the field evolves from θ to θ_f is

$$N(\theta) = \frac{8\pi}{3m_{\text{Pl}}^2} \int_{\theta_f}^{\theta} \frac{V(\theta)}{V'(\theta)} d\theta \tag{4.4}$$

where $V'(\theta) = \partial V/\partial\theta$ and θ_f is the value of the field at which inflation stops (reheating commences) and is obtained from

$$\epsilon(\theta_f) \equiv \frac{m_{\text{Pl}}^2}{16\pi} \left[\frac{V'(\theta_f)}{V(\theta_f)} \right]^2 = 1 \tag{4.5}$$

An upper limit on the initial value of the field is obtained by imposing $N(\theta_i) = 60$. Around this value, quantum fluctuations on scales observed today were produced and its size is given by [12]

$$\delta\rho/\rho \simeq \frac{(V(\theta_i))^{3/2}}{m_{\text{Pl}}^3 V'(\theta_i)} \tag{4.6}$$

The spectral index of these density perturbations and its dependence on the scale read

$$n_s - 1 = -6\epsilon(\theta_i) + 2\eta(\theta_i) \tag{4.7}$$

$$\frac{dn_s}{d\ln k} = -16\epsilon(\theta_i)\eta(\theta_i) + 24\epsilon(\theta_i)^2 + 2\xi^2(\theta_i) \tag{4.8}$$

where η and ξ^2 are the second and third slow-roll parameters

$$\eta = \frac{m_{\text{Pl}}^2 V''}{8\pi V} \quad \text{and} \quad \xi^2 = \frac{m_{\text{Pl}}^4 V'' V'}{(8\pi)^2 V^2}. \quad (4.9)$$

Expressed in terms of the parameters of the model, these quantities are given by

$$\sin(\theta_i) = [\beta(2-\beta)^2]^{1/2} \quad (4.10)$$

$$\delta\rho/\rho \simeq \frac{M^2 v}{m_{\text{Pl}}^3} \frac{(2-\beta)}{\beta^{1/2}} \quad (4.11)$$

$$\epsilon = \frac{m_{\text{Pl}}^2}{16\pi v^2} \frac{\beta}{(2-\beta)} \quad (4.12)$$

$$\eta = \frac{m_{\text{Pl}}^2}{8\pi v^2} \frac{(1-\beta)}{(2-\beta)} \quad (4.13)$$

$$n_s - 1 = \frac{m_{\text{Pl}}^2}{16\pi v^2} \frac{2(2+\beta)}{(2-\beta)} \quad (4.14)$$

$$\xi^2 = -\frac{m_{\text{Pl}}^4}{(8\pi)^2 v^4} \frac{\beta^{1/2}(1-\beta)}{(2-\beta)^{3/2}} \quad (4.15)$$

where

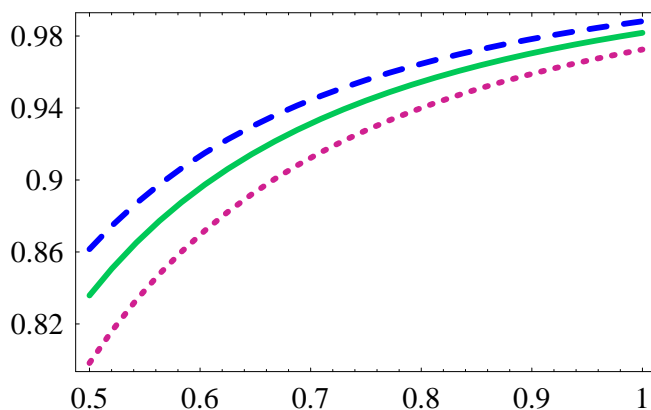
$$\beta = \frac{2 e^{-2Ny}}{y(1+y)} \quad \text{with} \quad y = \frac{m_{\text{Pl}}^2}{16\pi v^2} \quad (4.16)$$

and N number of e-folds before the end of inflation at which observable perturbations were generated.

These simplified expressions, although very easy to work with, hide the dependence of the cosmological observables on the model parameters. This becomes apparent in the fact that all the quantities above, but $\delta\rho/\rho$, depend exclusively on v . In order to link the original model parameters, v , g and g' to observations, the full potential must be used. Nevertheless, the slow-roll approximation is still reasonable.

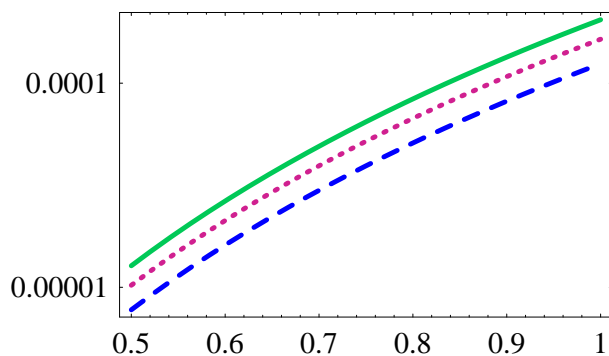
As can be seen from figure 1, the spectral index of the density fluctuations defines the range of values the parameter v can take. Clearly, although both coupling do contribute to the value of the spectral index, cosmological observations force v to live in the range $0.7 m_{\text{Pl}} < v < 0.9 m_{\text{Pl}}$ for any reasonable choice of g and g' . In this figure the spectral index is evaluated at a value θ_i such that sufficient inflation occurs when the field rolls down from θ_i to the end of inflation.

Once the symmetry breaking scale is defined to take values within this range, we can resort to the magnitude of density fluctuations to see whether some information on the couplings can be obtained. The answer is depicted on figure 2: the size of the primordial perturbations, evaluated at a value θ_i such that sufficient inflation occurs, does strongly depend on the couplings. However, it depends through the combination $(g^3 g')^{1/2}$, so only this combination can be bounded using the magnitude of density perturbations, and it must be $(g^3 g')^{1/2} \sim 10^{-5}$. It is important to notice that the scale v and the coupling



	g	g'
green (solid)	0.01	0.001
blue (dashed)	0.01	0.0001
purple (dotted)	0.1	0.03

Figure 1. Spectral index of the complete potential as a function of the symmetry breaking scale v for different values of the couplings.



	g	g'
green (solid)	0.01	0.001
blue (dashed)	0.01	0.0003
purple (dotted)	0.1	0.00001

Figure 2. Magnitude of the density fluctuations as a function of the symmetry breaking scale v for different values of the couplings

g determine not only the the mass of the inflaton (provided $g' \ll g$) but also the mass of the right-handed neutrino and the scale of spontaneous symmetry breaking; thus the importance of connecting both values with cosmological observations.

As an example, here is a set of input parameters from the cutoff scale that will produce

a satisfactory model of inflation:

$$\Lambda = m_0 = 10^{19} \text{ GeV} ; \quad G\Lambda^2 = 0.1 ; \quad \frac{8\pi^2}{N_f} = 0.05 . \quad (4.17)$$

From which we derive $\delta = 1$ and $g \sim 0.1$. We then require $g' \sim 10^{-7}$ to get the right magnitude of density perturbations. In this example the symmetry breaking scale v is close to the Planck scale, while m_R and M are of order 10^{18} GeV. Escaping from the large N_f limit assumed here would require a better understanding of the dynamics responsible for the condensate.

To bound the value of the couplings g and g' separately, a different observable with a different dependence on the couplings needs to be found. Two possibilities are at hand, although neither will be measured any time soon. Nevertheless, future measurements can rule out a dynamical origin of the inflaton field, as the one proposed here.

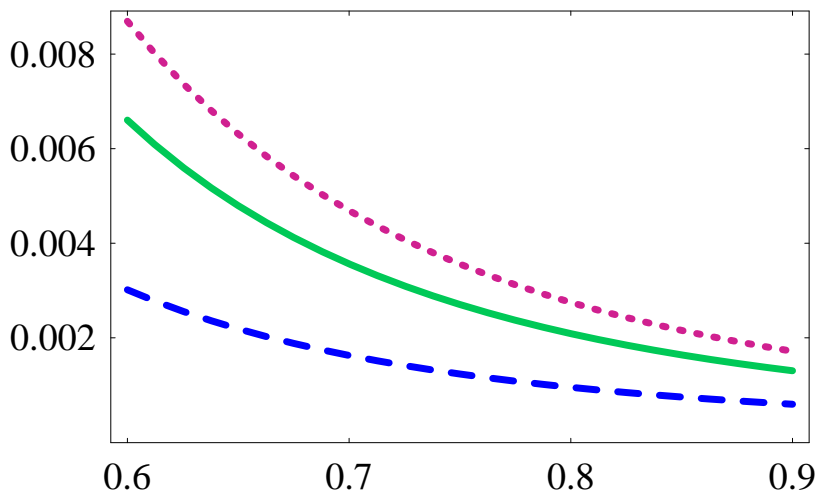
In addition to scalar (density) perturbations, our field will also give rise to tensor (gravitational wave) perturbations. Generally, the tensor amplitude is given in terms of the tensor/scalar ratio

$$r \equiv \frac{P_T}{P_R} = 16\epsilon \quad (4.18)$$

which is shown in Fig(3). The tensor to scalar ratio r goes like $g^2 g'^2$, and would offer the possibility of bounding each coupling individually, if the tensor amplitude were not well below the detection sensitivity of current and (near) future experiments, i.e. gravity waves are exponentially suppressed relative to the adiabatic scalar fluctuations over the observable large scale waveband. Gravity waves are the holy grail of next generation of experiments [13] and if found, will rule out this model.

Notice that in order to describe scalar and tensor fluctuations, only four parameters are needed (if we ignore the running): the amplitude and the spectra of both modes. The spectral indexes n_s and $n_T \equiv -2\epsilon$ characterize the latter, while the size of the scalar perturbations is basically characterized by the height of the potential (given by M^4 in the approximated expressions). The tensor amplitude is given by r . However, the tensor index is not an independent parameter since it is related to the tensor/scalar ratio by the inflationary consistency relation $r = -8n_T$ and therefore it is not useful for disentangling the values of each coupling.

In general, n_s is not a constant, and its dependence on the scale can be characterized by its running. Unfortunately the slow-roll approximation is numerically inaccurate for this parameter and may lead to discrepancies of a factor 2-3. However, we are interested in the order of magnitude of the result and therefore using the slow-roll approximation will leave our conclusions unaffected. As shown in Fig(4) our model predicts a very small and negative spectral index running, scaling as g'/g . It is so negligible small that it is essentially indistinguishable from zero running. Small scale CMB experiments [14] will provide more stringent tests on the running. If these experiments exclude a trivial (consistent with zero) running, i.e. if they detect a strong running, our model would be ruled out.



	g	g'
green (solid)	0.001	0.003
blue (dashed)	0.1	0.000003
purple (dotted)	0.01	0.005

Figure 3. Tensor to scalar ratio as a function of the symmetry breaking scale v for different values of the couplings

5 Conclusions

We have shown that the scalar field that drives inflation can have a dynamical origin, being a strongly coupled right handed neutrino condensate. The fact that Φ behaves like a sensible propagating field is a signal that we have chosen the correct low-energy degrees of freedom by introducing it. As the theory containing Φ is equivalent to a theory entirely written in terms of neutrino degrees of freedom, the field Φ can be interpreted as a right handed neutrino bound state.

The resulting model is phenomenologically tightly constrained, and can be experimentally (dis)probed in the near future. Probably the least attractive feature of the model is the range of values the symmetry breaking scale v is bounded to take, quite close to the Plank scale. This won't be the case in a scenario with more than one generation of right handed neutrinos. We wish to emphasize however, that the mass of the right handed neutrino (a crucial ingredient to obtain the right light neutrino spectrum within the see-saw mechanism in a complete three generation framework) is related to that of the inflaton and both completely determine the inflation features that can be tested by current and planned experiments. Thus, despite its problems we feel that the proposed dynamical origin of the inflaton field is sufficiently interesting to merit attention.

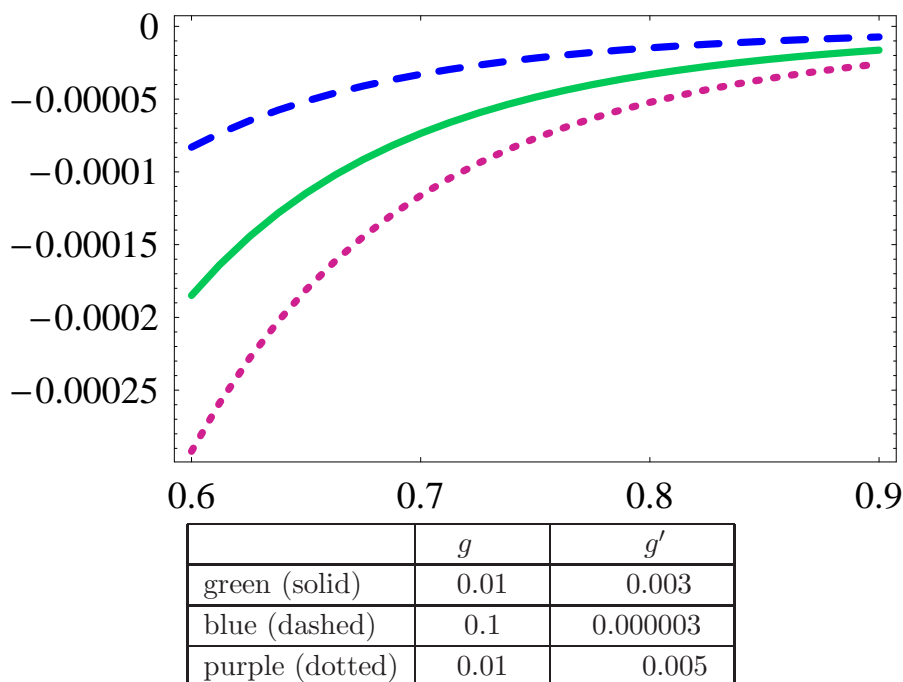


Figure 4. Running of the spectral index as a function of the symmetry breaking scale v for different values of the couplings

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References

- [1] A.H. Guth, *The inflationary universe: a possible solution to the horizon and flatness problems*, *Phys. Rev. D* **23** (1981) 347 [SPIRES];
A.D. Linde, *A new inflationary universe scenario: a possible solution of the horizon, flatness, homogeneity, isotropy and primordial monopole problems*, *Phys. Lett. B* **108** (1982) 389 [SPIRES];
A.J. Albrecht and P.J. Steinhardt, *Cosmology for Grand Unified Theories with radiatively induced symmetry breaking*, *Phys. Rev. Lett.* **48** (1982) 1220 [SPIRES].
- [2] Y. Nambu, *New theories in physics*, proceedings of the *XI International Symposium on Elementary Particle Physics*, Kazimierz, Poland, 1988, Z. Ajduk et al. eds., World Scientific, Singapore (1989).
- [3] W.A. Bardeen, C.T. Hill and M. Lindner, *Minimal dynamical symmetry breaking of the standard model*, *Phys. Rev. D* **41** (1990) 1647 [SPIRES].

- [4] A heavy, long-lived right handed neutrino can have a strong impact on the early universe, see G. Barenboim and O. Vives, *About a (standard model) universe dominated by the right matter*, *Phys. Rev. D* **79** (2009) 033007 [[arXiv:0806.4389](#)] [[SPIRES](#)].
- [5] K. Freese, J.A. Frieman and A.V. Olinto, *Natural inflation with pseudo-Nambu-Goldstone bosons*, *Phys. Rev. Lett.* **65** (1990) 3233 [[SPIRES](#)];
F.C. Adams, J.R. Bond, K. Freese, J.A. Frieman and A.V. Olinto, *Natural inflation: particle physics models, power law spectra for large scale structure and constraints from COBE*, *Phys. Rev. D* **47** (1993) 426 [[hep-ph/9207245](#)] [[SPIRES](#)].
- [6] E.W. Kolb and M.S. Turner, *The early universe*, *Front. Phys.* **69** (1990) 1 [[SPIRES](#)]
- [7] R. Holman et al., *Solutions to the strong CP problem in a world with gravity*, *Phys. Lett. B* **282** (1992) 132 [[hep-ph/9203206](#)] [[SPIRES](#)];
M. Kamionkowski and J. March-Russell, *Planck scale physics and the Peccei-Quinn mechanism*, *Phys. Lett. B* **282** (1992) 137 [[hep-th/9202003](#)] [[SPIRES](#)].
- [8] S.B. Giddings and A. Strominger, *Loss of incoherence and determination of coupling constants in quantum gravity*, *Nucl. Phys. B* **307** (1988) 854 [[SPIRES](#)].
- [9] G. Gilbert, *Wormhole induced proton decay*, *Nucl. Phys. B* **328** (1989) 159 [[SPIRES](#)];
S.-J. Rey, *The axion dynamics in wormhole background*, *Phys. Rev. D* **39** (1989) 3185 [[SPIRES](#)].
- [10] K. Freese, C. Savage and W.H. Kinney, *Natural inflation: the status after WMAP 3-year data*, *Int. J. Mod. Phys. D* **16** (2008) 2573 [[arXiv:0802.0227](#)] [[SPIRES](#)];
C. Savage, K. Freese and W.H. Kinney, *Natural inflation: status after WMAP 3-year data*, *Phys. Rev. D* **74** (2006) 123511 [[hep-ph/0609144](#)] [[SPIRES](#)].
- [11] WMAP collaboration, E. Komatsu et al., *Five-year Wilkinson Microwave Anisotropy Probe (WMAP) observations: cosmological interpretation*, *Astrophys. J. Suppl.* **180** (2009) 330 [[arXiv:0803.0547](#)] [[SPIRES](#)];
WMAP collaboration, D.N. Spergel et al., *Wilkinson Microwave Anisotropy Probe (WMAP) three year results: implications for cosmology*, *Astrophys. J. Suppl.* **170** (2007) 377 [[astro-ph/0603449](#)] [[SPIRES](#)].
- [12] J.E. Lidsey et al., *Reconstructing the inflaton potential: an overview*, *Rev. Mod. Phys.* **69** (1997) 373 [[astro-ph/9508078](#)] [[SPIRES](#)].
- [13] See e.g. FTQ collaboration, D. Samtleben et al., *QUIET - Measuring the CMB polarization with coherent detector arrays*, [arXiv:0806.4334](#) [[SPIRES](#)];
PLANCK collaboration, F.R. Bouchet, *The Planck satellite: status & perspectives*, *Mod. Phys. Lett. A* **22** (2007) 1857 [[SPIRES](#)];
J. Bock et al., *Task force on cosmic microwave background research*, [astro-ph/0604101](#) [[SPIRES](#)].
- [14] C. Dickinson et al., *High sensitivity measurements of the CMB power spectrum with the extended very small array*, *Mon. Not. Roy. Astron. Soc.* **353** (2004) 732 [[astro-ph/0402498](#)] [[SPIRES](#)];
B.S. Mason et al., *The anisotropy of the microwave background to $l = 3500$: deep field observations with the cosmic background imager*, *Astrophys. J.* **591** (2003) 540 [[astro-ph/0205384](#)] [[SPIRES](#)];
ACBAR collaboration, C.-L. Kuo et al., *High resolution observations of the CMB power spectrum with ACBAR*, *Astrophys. J.* **600** (2004) 32 [[astro-ph/0212289](#)] [[SPIRES](#)].