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Sneutrino chaotic inflation and landscape

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ABSTRACT

The most naive interpretation of the BICEP2 data is the chaotic inflation by an inflaton with a quadratic potential. When combined with supersymmetry, we argue that the inflaton plays the role of right-handed scalar neutrino based on rather general considerations. The framework suggests that the right-handed sneutrino tunneled from a false vacuum in a landscape to our vacuum with a small negative curvature and suppressed scalar perturbations at large scales.

1. Introduction

Discoveries of the B-mode polarization by BICEP2 [1] and the Higgs boson by ATLAS [2] and CMS [3] mark a huge progress in fundamental physics. BICEP2 result, if confirmed, suggests an inflaton potential with a large field amplitude, where a simple quadratic potential $V = \frac{1}{2}M^2\phi^2$ [4] is preferred. Because the potential needs to maintain this form up to an amplitude of 15–16 times the reduced Planck scale, the inflaton field most likely does not participate in gauge interactions to avoid large radiative corrections. On the other hand, the implied mass scale $M \simeq 2 \times 10^{13}$ GeV [5] is much larger than the observed Higgs mass of 126 GeV, hinting at a mechanism to protect a large hierarchy, such as supersymmetry.

Once inflation is considered proven, the immediate next question is what the inflaton is. In particular, we need to know how inflation ends and reheats the Universe, and how the baryon asymmetry is created after inflation, given that inflation wipes out any pre-existing baryon asymmetry. On both questions, it is clearly important to know how the inflaton couples to the known particles in the Standard Model.

We show in this Letter that, being a gauge singlet, the inflaton naturally induces the neutrino mass. Therefore, it is possible to identify the inflaton with the right-handed scalar neutrino [6]. Then leptogenesis [7] is the likely mechanism for creating the baryon asymmetry. Tantalizingly, the picture is suggestive of the decay of false vacuum in the landscape [8], where the right-handed scalar neutrino tunnels from a local minimum to our minimum. If so, suppression of scalar perturbation at low ℓ [9] and a small negative curvature are expected.

2. *B*-mode

Cosmic inflation was originally proposed to solve the flatness and horizon problems of the big bang cosmology [10,11].[1] The graceful exit problem of the original inflation was solved by the new inflation [15,16] where the slow-roll inflaton drives the exponential expansion. At the same time, it became the dominant paradigm to generate the nearly scale-invariant, adiabatic, and Gaussian density perturbations from the quantum fluctuation of the inflaton field. Its prediction has been known to explain the data very well, including anisotropy in cosmic microwave background radiation (CMB) and galaxy power spectrum. However, the inflation paradigm so far lacked the definitive proof.

Primordial B-mode polarization of CMB is regarded as a possible definite proof of inflation [17–19]. If the expansion rate is very high during the inflationary period, gravitons are created due to the quantum fluctuation. Once the mode exits the horizon,

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[1] The exponentially expanding universe was also studied in Refs. [12–14], where the flatness and horizon problems were not discussed.

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the quantum fluctuation becomes classical and the gravitons are imprinted as primordial gravitational waves, i.e., tensor perturbations of the space–time metric [20–22]. The CMB photons acquire polarization through Thomson scatterings with electrons on the last scattering surface because of local quadrupole anisotropies at each point. While density (scalar) perturbations induce only E-mode polarization, tensor perturbations induce both E-mode and B-mode polarization patterns. Most importantly, B-mode polarization at small multipoles is unlikely to be generated by other mechanisms. The B-mode polarization at large multipoles, on the other hand, is induced by gravitational lensing effect of large-scale structures such as clusters of galaxies on the way from the last scattering surface to the Earth.[2]

The recent data from BICEP2 experiment [1] may have provided such a long-awaited proof of the inflation paradigm. It reported a detection of the B-mode polarization, which can be explained by the tensor-to-scalar ratio, $r = 0.20^{+0.07}_{-0.05}$. Taken at face value, the BICEP2 results exclude many inflation models and strongly suggest large-field inflation models in which the inflaton field amplitude during inflation exceeds the reduced Planck mass, $M_{Pl} = G_N^{-1/2}/\sqrt{8\pi} \simeq 2.4 \times 10^{18}$ GeV. Among various large-field inflation models, by far the simplest and therefore most attractive one is the chaotic inflation with a simple quadratic potential [4], which predicts $r \simeq 0.13(0.16)$ and $n_s \simeq 0.97(0.96)$ for the e-folding number $N_e = 60(50)$, completely consistent with the data. The inflaton mass is fixed to be $M \simeq 2 \times 10^{13}$ GeV by the normalization of density perturbations [5].

It is worth noting that the BICEP2 results are in tension with the Planck data [24] on the relative size of density perturbations on large and small scales. This tension could be due to some unusual features in the density perturbations such as a negative running of the scalar spectral index [1,25,26], or it may indicate the decay of a false vacuum in a landscape just before the beginning of slow-roll inflation [8,9]. The apparent tension between BICEP2 and Planck will also be partially relaxed if the true value of the tensor-to-scalar ratio is close to the lower end of the observed range. We will return to this issue later in this Letter.

3. Inflaton

Given the inflation potentially proven, now the community should move on to a new question: what is the inflaton? This is a pressing question in order to understand the subsequent cosmic history after inflation, which depends on the coupling of the inflaton to the standard-model particles.

On the other hand, the large hierarchy between the inflaton mass and the Higgs mass needs to be protected by supersymmetry (see [27] for a review). Even though the LHC has not discovered superparticles yet, it can be hidden due to a degenerate spectrum (see [28] for a mechanism to create a degenerate spectrum) or is simply somewhat heavier than anticipated. The minute supersymmetry is introduced, we need a matter parity to avoid too-fast proton decay. The matter parity of the inflaton field can be either even or odd. This is a crucial question in order to see how the inflaton couples to the standard-model particles.

If the inflaton field Φ has an even matter parity, its lowest-order coupling to the standard-model particles is $W = \lambda \Phi H_u H_d$. Inflaton decays into Higgs fields and reheats the Universe.[3] In this case, we do not see any obvious connection of the inflaton properties to low-energy observables, nor to the origin of baryon asymmetry.

On the other hand, if the inflaton has an odd matter parity, the lowest-order coupling of the inflaton to the standard-model particles is $W = h\Phi L H_u$, where L and H_u represent the lepton doublet and up-type Higgs superfields, respectively, and the flavor index is suppressed. We expect the low-energy consequence to be $(LH_u)^2/M$, which is nothing but the neutrino mass. In other words, we may say small neutrino mass is a low-energy consequence of the inflaton.

In fact, the suggested mass of the inflaton is very close to that of the right-handed neutrino $\approx 10^{14}$ GeV required in the seesaw mechanism that explains small neutrino masses [36–41]. It is indeed a gauge-singlet. Then we can identify the inflaton with the right-handed scalar neutrino, as proposed some time ago [6].

To make the discussion more concrete, let us pick a simple model of chaotic inflation by a quadratic potential within supergravity [42,43]. The superpotential is simply the mass term

$$W = MX\Phi, \tag{1}$$

while the Kähler potential has a shift symmetry for the field $\Phi \to \Phi + ic$,

$$K = \frac{1}{2}\left(\Phi^* + \Phi\right)^2 + X^* X + \text{higher orders}. \tag{2}$$

We assign the odd matter parity to Φ and X.

The scalar potential in supergravity reads

$$V = e^{K/M_{Pl}^2}\left(K^{I\bar{J}}(D_I W)(D_J W)^\dagger - 3\frac{|W|^2}{M_{Pl}^2}\right), \tag{3}$$

with $D_I W = \partial_I W + (\partial_J K)W/M_{Pl}^2$ and the subscript I is a label for a scalar field.

Most importantly, the imaginary component of Φ does not appear in the Kähler potential because of the shift symmetry, and therefore the potential along Im[Φ] remains relatively flat at super-Planckian field values.

Let us first suppose that, during inflation, all the other fields are stabilized at the origin. It is then straight-forward to work out the potential for X and Φ,

$$V = e^{K/M_{Pl}^2}\left(\left|\frac{(\Phi^* + \Phi)}{M_{Pl}^2} MX\Phi + MX\right|^2 + \left|\frac{X^*}{M_{Pl}^2} MX\Phi + M\Phi\right|^2 - 3\left|\frac{M}{M_{Pl}} X\Phi\right|^2\right). \tag{4}$$

Specializing to the imaginary term $\Phi = i\phi/\sqrt{2}$ and $X = 0$, we find

$$V = \frac{1}{2}M^2\phi^2, \tag{5}$$

a simple quadratic potential, and the correct size of density perturbations is generated for $M \simeq 2 \times 10^{13}$ GeV. X has the same mass for the above Kähler potential, but it can be stabilized at the origin with a positive mass of order the Hubble parameter, by adding a quartic coupling $\delta K = -|X|^4$ in the Kähler potential.

Note that the lowest-order couplings of Φ and X to the standard-model fields allowed by the odd matter parity is

$$W_{\text{coupl}} = h_\alpha \Phi L_\alpha H_u + \tilde{h}_\alpha X L_\alpha H_u. \tag{6}$$

Then the low-energy consequence is indeed the neutrino mass

$$W_{\text{eff}} = \frac{1}{M}(h_\alpha L_\alpha H_u)(\tilde{h}_\beta L_\beta H_u). \tag{7}$$

[2] The lensing B-mode was recently measured by the POLARBEAR experiment [23].

[3] In this case, the inflaton can decay into a pair of gravitinos because of a possible Kähler potential term linear in the inflaton field [29–34], and tight cosmological constraints were obtained [35]. See later discussions in this Letter how the constraints can be evaded.

Therefore, it is tempting to assume that the inflaton takes part in the origin of neutrino mass, and we hereafter identify the inflaton field Φ as one of the right-handed neutrinos in the seesaw mechanism.[4]

4. Three right-handed neutrinos

To be explicit, let us consider the case of three right-handed neutrinos, although the number of right-handed neutrinos is not restricted, e.g., by anomaly cancellations. The following arguments can be straightforwardly applied to a case with right-handed neutrinos different from three.[5]

The superpotential for the right-handed neutrinos N_i is

$$W = \frac{1}{2} M_{ij} N_i N_j + h_{i\alpha} N_i L_\alpha H_u, \tag{8}$$

where M_{ij} $(i, j = 1, 2, 3)$ is the mass matrix for the right-handed neutrinos and $h_{i\alpha}$ $(\alpha = e, \mu, \tau)$ denotes the Yukawa coupling of the right-handed neutrino with the lepton doublet L_α and the up-type Higgs H_u. Integrating out the heavy right-handed neutrinos, one obtains the seesaw formula for the light neutrino mass [36–41],

$$(m_\nu)_{\alpha\beta} = h_{i\alpha} \left(M^{-1}\right)_{ij} h_{j\beta} v_u^2, \tag{9}$$

where $v_u = \langle H_u \rangle$ is the expectation value of the Higgs field.

For successful inflation, we assume

$$M_{ij} = \begin{pmatrix} m & 0 & 0 \\ 0 & 0 & M \\ 0 & M & 0 \end{pmatrix} \tag{10}$$

while the Kähler potential respects a shift symmetry for N_3,

$$K = N_1^\dagger N_1 + N_2^\dagger N_2 + \frac{1}{2}\left(N_3^\dagger + N_3\right)^2 + \cdots, \tag{11}$$

where the dots represent higher order terms, and we suppressed L_i and H_u as they can be stabilized at the origin during inflation. It is the imaginary component of N_3 that becomes the inflaton, and the inflaton potential is given by (5). This is a simple but new realization of the right-handed sneutrino inflation based on the chaotic inflation model (1).

5. Reheating and leptogenesis

In order to generate a sufficiently large neutrino mass $m_\nu \approx 0.05$ eV, we need the Yukawa coupling as large as $h \sim 0.1$, where h denotes the typical value of $h_{i\alpha}$. Then the inflaton reheats the Universe up to

$$T_{RH} \approx g_*^{-\frac{1}{4}} \sqrt{\frac{h^2}{8\pi} M M_{Pl}} \gtrsim 10^{13}\ \text{GeV}, \tag{12}$$

where g_* counts the relativistic degrees of freedom in thermal plasma. For such high reheating temperature, the e-folding number N_e is about 60, and the predicted values of r and n_s are $r \simeq 0.13$ and $n_s \simeq 0.97$. Also, the right-handed neutrinos thermalize after reheating and the usual thermal leptogenesis takes place.[6] By keeping $m \ll M$, the N_1 plays the dominant role in leptogenesis. The CP asymmetry in its decay is given by (see, e.g., [45])

$$\epsilon_1 = \frac{1}{4\pi} \frac{\Im m \sum_{\alpha,\beta} (h_{1\alpha} h_{1\beta} h_{2\alpha}^* h_{3\beta}^*)}{\sum_\alpha h_{1\alpha} h_{1\alpha}^*} \frac{m}{M}. \tag{13}$$

$m > 4 \times 10^8$ GeV is required for a successful thermal leptogenesis [46].

Given the high reheating temperature Eq. (12), the gravitino is copiously produced by thermal scatterings, and its abundance is given by [47–51],

$$Y_{3/2} \simeq 2 \times 10^{-9} \left(\frac{T_{RH}}{10^{13}\ \text{GeV}}\right), \tag{14}$$

where $Y_{3/2}$ is the ratio of the gravitino number density to the entropy density, and we suppressed the contributions from longitudinal mode. Non-thermal gravitino production from inflaton decays is absent due to the matter parity [29]. The gravitino decays into the Lightest Supersymmetric Particle (LSP) at a later time. If LSP is stable and weighs about TeV, it exceeds the observed dark matter abundance by about four orders of magnitude [47]. Therefore, one possible solution is the LSP to be as light as 100 MeV.

This problem can be avoided also in the following ways. One possibility is that the gravitino is light [52,53], such as in low-scale gauge mediation models [54]. The light gravitinos are thermalized and account for the observed dark matter abundance in the presence of mild entropy production by e.g. the lightest messengers [55,56]. In particular, if the gravitino mass is lighter than about 16 eV [57], there is no cosmological bound on gravitinos, as their contribution to (hot) dark matter becomes negligibly small. If this is the case, the NLSP would decay inside the collider detector, and the gravitino mass can be determined in the future experiments by measuring the branching fractions of two decay modes of NLSP [58,59].

Another possibility is to allow for a small matter parity violation. Then, the LSPs are unstable and decay before big bang nucleosynthesis. Non-thermally produced LSPs via the gravitino decay will also decay before big bang nucleosynthesis, if the gravitino mass is heavier than a few tens of TeV. The matter parity violation is bounded above in order not to erase the baryon asymmetry generated by leptogenesis [60]. For such small matter parity violation, the LSPs decay at macroscopic distances from the beam line at collider experiments. If the LSP is electrically charged as in the case of the stau LSP, it gives rise to highly ionizing tracks (see e.g. Refs. [61–63]). If the matter parity violation is sufficiently large (but still below the upper bound for successful leptogenesis), the LSPs may decay within the collider detector, leading to kinked charged tracks [64] or displaced vertices [65].

If there are many singlets with an odd matter parity around the inflaton mass scale, they also contribute to the neutrino masses, and so, the contributions of the inflaton can be subdominant. If its couplings are as small as $h \sim 10^{-5}$, the reheating temperature will be of order 10^9 GeV, greatly relaxing the gravitino overproduction problem. In particular, the gravitino mass about 100–1000 TeV preferred by the anomaly mediation [66,67], the pure gravity mediation scenario [68,69], or the minimal split SUSY [70,71] is allowed even without the matter parity violation. Interestingly, the observed Higgs mass of 126 GeV can be naturally explained in this case.

6. Landscape

The shift symmetry on N_3 is only approximate as it is explicitly broken by the neutrino Yukawa coupling $h \sim 0.1$. Regarding h as a spurion, we expect corrections of the form

$$V = \frac{1}{2} M^2 \phi^2 \left(1 + c \frac{h^2 \phi^2}{M_{Pl}^2} + \cdots\right), \tag{15}$$

[4] In principle, this contribution can saturate the whole neutrino mass matrix with two non-zero eigenvalues.

[5] In F-theory compactifications, complex structure moduli can be identified with right-handed neutrinos and their number is typically much lager than three [44].

[6] Thermal leptogenesis takes place even when the reheating proceeds efficiently through preheating and the subsequent dissipation processes.

with $c = \mathcal{O}(1)$. Namely, the potential deviates from the simple quadratic potential at the amplitude $\phi \sim M_{Pl}/h$. This is approximately the amplitude that corresponds to the e-folding $N \sim 60$ to solve the flatness and horizon problems of the Universe. It suggests the possibility that the inflation was "just so": the total e-folding number was just about 60.

It is interesting to note that the just-so e-folding of $N \sim 60$ is what is expected in the landscape [72]. If there are a large number of local minima in the potential, the tunneling from the local minimum closest to "our" minimum should set off the inflation. Since a flat potential required for inflation is not generic, the anthropic argument suggests the e-folding is as small as necessary for us to exist. Our existence requires a low enough curvature to allow for a successful structure formation, that leads to the lower limit $N \gtrsim 60$. It corresponds to the initial amplitude $\phi \sim 15 M_{Pl}$.

If taken seriously, the overall picture suggests that the right-handed scalar neutrino tunneled from a local minimum to our minimum by the Coleman–De Luccia mechanism. This is a truly remarkable role for the neutrino. As discussed in Refs. [72,73], we then expect that the tunneling from the local minimum brings the right-handed sneutrino where the potential is steeper than ϕ^2, and the field starts to roll down the potential. Yet the field rolls very slowly because the large negative curvature required in the Coleman–De Luccia mechanism acts as a friction in the field equation for the right-handed sneutrino, solving the overshoot problem. At the same time, the right-handed sneutrino is homogenized over many horizons solving the initial condition problem for the chaotic inflation. Only after the curvature is sufficiently flattened out, the field starts to roll faster. Given the steeper potential at the beginning, it results in a suppression in scalar perturbation at low ℓ due to its faster motion $\delta\rho/\rho \propto V'/\dot{\phi}^2$, which ameliorates the tension with the Planck data on the temperature anisotropy [8,9]. A small negative curvature will remain at the level of $\Omega_k \sim 10^{-4}$–10^{-2}. We look forward to future precise measurements on Ω_k from large-scale deep spectroscopic surveys such as SuMIRe [74].

7. Conclusion

In this Letter, we have argued that the inflaton of chaotic inflation with a quadratic potential suggested by the BICEP2 data can be naturally identified with a right-handed scalar neutrino. This explains why the suggested mass of the inflaton is very close to that of the right-handed neutrino $\approx 10^{14}$ GeV required in the seesaw mechanism. The size of the Yukawa coupling to generate neutrino mass violates the shift symmetry, making the total e-folding "just so", i.e., $N \simeq 60$. This is what is expected in the landscape. The leptogenesis takes place as the reheating process itself, or thermally after reheating.

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