

Updating GUT-scale pole Higgs inflation after ACT DR6

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We consider models of chaotic inflation driven by the real parts of a conjugate pair of Higgs superfields involved in the spontaneous breaking of a grand unification symmetry at a scale assuming its value within the minimal supersymmetric Standard Model (MSSM). We combine a superpotential, which is uniquely determined by applying a continuous R symmetry, with two fractional shift-symmetric Kähler potentials introducing two free parameters (p, N) . The inflationary observables provide an excellent match to the recent Atacama Cosmology Telescope data for $1.355 \leq p \leq 6.7$ and $6 \times 10^{-5} \leq N \leq 0.7$. The attainment of inflation allows for sub-Planckian inflaton values and possibly detectable primordial gravitational waves with (p, N) values of order unity. A solution to the μ problem of MSSM and baryogenesis via nonthermal leptogenesis can be also accommodated by embedding the model into a $B - L$ extension of the MSSM.

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I. INTRODUCTION

The announcement of data release 6 from the Atacama Cosmology Telescope (ACT) [1,2] fueled a plethora of works [3–42] that revise many well-motivated inflationary models—for a review, see Ref. [43]. This is because the combination of the aforementioned data with the cosmic microwave background measurements by Planck [44,45], together with the Dark Energy Spectroscopic Instrument baryon acoustic oscillation results [46] suggests a value of the (scalar) spectral index n_s substantially larger than the one indicated from the Planck data [44]. Namely, taking into account also the tensor-to-scalar ratio r , the so-called P-ACT-LB-BK18 data entails [2]

$$n_s = 0.9743 \pm 0.0068 \quad \text{and} \quad r \leq 0.038 \quad (1)$$

at 95% confidence level (CL).

In one recent paper [47], we propose a variant of pole inflation [48–53] that ensures compatibility with the (n_s, r) values in Eq. (1). This goal becomes possible by introducing a particular class of rational Kähler potentials K that include one extra parameter p compared to the α - (N in our notation) attractor models [54–56]. Most notably, this achievement is independent of the type of the monomial superpotential W selected leading, thereby, to the

establishment of a new set of (N, p) attractors. Working in the context of supergravity (SUGRA), in our present investigation we adapt our proposal in Ref. [47] to a grand unified theory (GUT) including the inflaton in a GUT-scale Higgs sector. In a such case, both W and K are consistent with a gauge symmetry rendering our scheme much more well defined. Following the terminology of Ref. [47], we here realize T_p -model inflation identifying the inflaton with the real part of a conjugate pair of Higgs superfields—cf. Refs. [52,53]. We therefore could name our model T_p -model Higgs inflation (HI) or, for short, p HI—for other works on GUT-scale HI, see, e.g., Refs. [57–65].

Apart from the compatibility with Eq. (1), p HI offers high enough r values that can be hopefully tested in the near future [66]. Nonetheless, it also allows us to explore consequences beyond the inflationary phenomenology. Namely, embedding p HI into a $B - L$ extension of the minimal supersymmetric Standard Model (MSSM), we can accommodate a resolution of the μ problem, following the mechanism in Refs. [53,67], and acceptable baryogenesis via nonthermal leptogenesis (nTL) [68–70] consistent with the gravitino ($\tilde{G}$) constraint [71,72]. The success of our postinflationary scenario bounds from above the ratio $\mu/m_{3/2}$ —where $m_{3/2}$ is the $\tilde{G}$ mass—and so it gives us the opportunity to link the high- to low-scale supersymmetry (SUSY).

In Sec. II, we describe the building blocks (W and K 's) of our proposal, we outline the derivation of the inflationary potential in Sec. III, and test our model against observations in Sec. IV. Finally, we expose a postinflationary completion in Sec. V and summarize our conclusions in Sec. VI. In the

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Appendix, we present our approach to the reheating stage in our model. Unless otherwise stated, we use units where the reduced Planck scale $m_P = 2.43 \times 10^{18}$ GeV is equal to unity.

II. MODEL SETUP

Our starting point is the superpotential known from the models of F-term hybrid inflation [37–41]

$$W = \lambda S(\tilde{\Phi}\Phi - M^2/2), \quad (2)$$

which is uniquely determined, at renormalizable level, by a gauge symmetry $\mathbb{G}$ and a convenient continuous R symmetry. Here, λ and M are two constants that can both be taken positive by field redefinitions; S is a left-handed superfield singlet under $\mathbb{G}$; $\tilde{\Phi}$ and Φ are a pair of left-handed superfields belonging to nontrivial conjugate representations of $\mathbb{G}$ and reducing its rank by their vacuum expectation values (VEVs). Just for definiteness, we restrict ourselves to $\mathbb{G}_{B-L} = \mathbb{G}_{\text{SM}} \times U(1)_{B-L}$ which consists of the simplest GUT beyond the MSSM—here $\mathbb{G}_{\text{SM}}$ is the gauge group of the MSSM and B and L denote the baryon and lepton number. The $B-L$ and R charges of S , $\tilde{\Phi}$, and Φ can be assigned as shown in Table 1 of Ref. [63]. At the SUSY limit, W leads to a $B-L$ phase transition, since, at the SUSY vacuum,

$$\langle S \rangle \simeq 0 \quad \text{and} \quad \langle \Phi \rangle = \langle \tilde{\Phi} \rangle = M/\sqrt{2} \quad \text{for } M \ll 1. \quad (3)$$

$U(1)_{B-L}$ is spontaneously broken via $\langle \Phi \rangle$ and $\langle \tilde{\Phi} \rangle$.

To realize $p\text{HI}$ using W in Eq. (2)—with S stabilized at the origin and the inflaton included in the $\tilde{\Phi}$ – Φ system—a careful choice of K is imperative as in Refs. [52,62,63]. The K proposed here includes two contributions without mixing between $\tilde{\Phi}$ – Φ and S , i.e.,

$$\tilde{K}_{\ell\text{st}} = \tilde{K}_{\ell} + K_{\text{st}} \quad \text{with} \quad \ell = 1, 2 \quad (4)$$

from which K_{st} successfully stabilizes S along the inflationary path—see below—without invoking higher order terms. We adopt the form [73]

$$K_{\text{st}} = N_{\text{st}} \ln(1 + |S|^2/N_{\text{st}}) \quad \text{with} \quad 0 < N_{\text{st}} < 6, \quad (5)$$

which parametrizes the compact manifold $SU(2)/U(1)$ with curvature $2/N_{\text{st}}$. On the other hand, $\tilde{K}_{\ell}$ is devoted to the inflationary part and includes two contributions: a real one K_{ℓ} that determines the Kähler metric for the inflaton sector and an holomorphic (and antiholomorphic) part K_{sh} that provides $\tilde{K}_{\ell}$ with a shift symmetry assuring $\tilde{K}_{\ell} = 0$ during $p\text{HI}$ —cf. Ref. [47]. Therefore, $\tilde{K}_{\ell}$ has the structure

$$\tilde{K}_{\ell} = K_{\ell} + K_{\text{sh}} \quad \text{with} \quad \ell = 1, 2. \quad (6)$$

The key ingredient of our proposal is the selection of a totally fractional K_{ℓ} for $\ell = 1$, which reads

$$K_1 = \frac{N}{(1 - |\Phi|^2 - |\tilde{\Phi}|^2)^p} \quad (7a)$$

defined for $N > 0$ and $|\Phi|^2 + |\tilde{\Phi}|^2 < 1$. Identical results can be achieved if we select for $\ell = 2$

$$K_2 = \frac{N}{2(1 - 2|\Phi|^2)^p} + \frac{N}{2(1 - 2|\tilde{\Phi}|^2)^p} \quad (7b)$$

defined for $N > 0$, $|\Phi| < 1/\sqrt{2}$, and $|\tilde{\Phi}| < 1/\sqrt{2}$. Although K_2 seems to be more contrived than K_1 —due to the appearance of the same factor in the numerators and exponent in the denominators—both K 's share the same symmetries: These are invariant under $\mathbb{G}$ and the interchange $\Phi \leftrightarrow \tilde{\Phi}$. As regards the parameters (p, N) , we expect that N is constrained observationally—as for the N attractors [52–54]—and confine ourselves conservatively in the range $0.1 \leq p \leq 10$.

On the other hand, K_{sh} is common for both $\tilde{K}_{\ell}$ and takes the form

$$K_{\text{sh}} = -\frac{N}{2(1 - 2\tilde{\Phi}\Phi)^p} - \frac{N}{2(1 - 2\tilde{\Phi}^*\Phi^*)^p}, \quad (8)$$

where the asterisk denotes complex conjugation. It is obvious that K_{sh} let unchanged the Kähler metric by construction. Actually, the function in the denominators of K_{sh} is the unique possible quadratic function consistent with the gauge symmetry—cf. Eq. (2). Both K 's in Eq. (6) parametrize hyperbolic Kähler manifolds but without constant curvatures, unlike the case of T-model HI [52].

III. INFLATIONARY POTENTIAL

The appropriateness of the proposed W and $\tilde{K}_{\ell\text{st}}$ in Eqs. (2) and (4) can be verified if we identify the inflationary potential and the canonically normalized inflaton with those introduced in Ref. [47] for T_p -model inflation.

To embark on it, we focus on the SUGRA Lagrangian density for the complex scalar fields $z^{\alpha} = S, \Phi, \tilde{\Phi}$ —denoted by the same superfield symbol—which can be written as [74]

$$\mathcal{L} = \sqrt{-g} \left(-\frac{1}{2} \mathcal{R} + K_{\alpha\bar{\beta}} g^{\mu\nu} D_{\mu} z^{\alpha} D_{\nu} z^{*\bar{\beta}} - V \right), \quad (9)$$

where g is the determinant of the background Friedmann-Robertson-Walker metric $g^{\mu\nu}$ with signature $(+, -, -, -)$, and $\mathcal{R}$ is the Ricci scalar. The Kähler metric $K_{\alpha\bar{\beta}}$ and its inverse $K^{\bar{\beta}\alpha}$ are defined as follows:

$$K_{\alpha\bar{\beta}} = K_{,z^\alpha z^{\bar{\beta}}} > 0 \quad \text{with} \quad K^{\bar{\beta}\alpha} K_{\alpha\bar{\gamma}} = \delta_{\bar{\gamma}}^{\bar{\beta}}, \quad (10)$$

where z^α are complex scalar fields—the symbol, z^α as subscript denotes derivation with respect to z^α —and D_μ is the gauge covariant derivative. Also V is the SUGRA potential that can be found in terms of W and $\tilde{K}_{\ell\text{st}}$ in Eqs. (2) and (4) via the formula

$$V = e^K (K^{\alpha\bar{\beta}} F_\alpha F_{\bar{\beta}}^* - 3|W|^2) + \frac{g^2}{2} \sum_a D_a D_a, \quad (11a)$$

where the F- and D-terms are defined as follows:

$$F_\alpha = W_{,z^\alpha} + K_{,z^\alpha} W \quad \text{and} \quad D_a = z^\alpha (T_a)_\alpha^\beta K_\beta \quad (11b)$$

with $K_\alpha = K_{,z^\alpha}$ and the summation is applied over the generators T_a of $\mathbb{G}$. If we parametrize $\Phi, \bar{\Phi}$, and S as follows:

$$\Phi = \phi e^{i\theta} \cos \theta_\Phi, \quad \bar{\Phi} = \phi e^{i\bar{\theta}} \sin \theta_\Phi,$$

$$\text{and} \quad S = (s + i\bar{s})/\sqrt{2}, \quad (12)$$

with $0 \leq \theta_\Phi \leq \pi/2$, we can easily deduce from Eq. (11) that a D-flat direction occurs at

$$\langle \bar{s} \rangle_I = \langle s \rangle_I = \langle \theta \rangle_I = \langle \bar{\theta} \rangle_I = 0 \quad \text{and} \quad \langle \theta_\Phi \rangle_I = \pi/4, \quad (13)$$

where the symbol $\langle Q \rangle_I$ means the value of the quantity Q during $p\text{HI}$. Since $\langle K \rangle_I = 0$ and $\langle K_{SS^*} \rangle_I = 1$ for both $\tilde{K}_{\ell\text{st}}$ in Eq. (4) the only surviving term from V in Eq. (11a) is

$$V_{\text{HI}} = \langle V \rangle_I = \langle e^K K^{SS^*} |W_{,S}|^2 \rangle_I = \lambda^2 (\phi^2 - M^2)^2/4, \quad (14)$$

which reduces to the quartic potential considered in Ref. [47] for $M \ll 1$ (see below).

To specify the canonically normalized inflaton, we note that, for both $\tilde{K}_{\ell\text{st}}$ in Eq. (4), $K_{\alpha\bar{\beta}}$ along the configuration in Eq. (13) takes the form

$$\langle K_{\alpha\bar{\beta}} \rangle_I = \text{diag}(M_{\Phi\Phi}, \langle K_{SS^*} \rangle_I), \quad (15)$$

where the matrix $M_{\Phi\Phi}$ is specified as follows:

$$M_{\Phi\Phi} = \frac{pN}{2f_{\text{T}}^{p+2}} \begin{cases} \begin{pmatrix} f_{\text{T}} + f_p & (p+1)\phi^2 \\ (p+1)\phi^2 & f_{\text{T}} + f_p \end{pmatrix} & \text{for } K = \tilde{K}_{1\text{st}}, \\ 2f_p \text{diag}(1, 1) & \text{for } K = \tilde{K}_{2\text{st}}, \end{cases} \quad (16)$$

where $f_{\text{T}} = 1 - \phi^2$ and $f_p = 1 + p\phi^2$. Upon diagonalization of $M_{\Phi\Phi}$ for $K = \tilde{K}_{1\text{st}}$ we find its eigenvalues, which are

TABLE I. Mass-squared spectrum for $K = \tilde{K}_{\ell\text{st}}$ ($\ell = 1, 2$) along the path in Eq. (13).

Fields	Eigenstates	Masses squared	
		$K = \tilde{K}_{1\text{st}}$	$K = \tilde{K}_{2\text{st}}$
Four real scalars	$\hat{\theta}_+$	$m_{\hat{\theta}_+}^2$	$6H_{\text{HI}}^2$
	$\hat{\theta}_\Phi$	$m_{\hat{\theta}_\Phi}^2$	$M_{BL}^2 + 6H_{\text{HI}}^2$
	$s, \bar{s}$	m_s^2	$6(1/N_{\text{st}} + 2f_{\text{T}}^{p+2}/pN\phi^2 f_p)H_{\text{HI}}^2$
One gauge boson	A_{BL}	M_{BL}^2	$2g^2 pN\phi^2/f_{\text{T}}^{p+1} \quad 2g^2 pN\phi^2 f_p/f_{\text{T}}^{p+2}$
Four Weyl spinors	$\hat{\psi}_\pm$	$m_{\hat{\psi}_\pm}^2$	$6f_{\text{T}}^{p+2}H_{\text{HI}}^2/pN\phi^2 f_p$
	$\lambda_{BL}, \hat{\psi}_{\Phi-}$	M_{BL}^2	$2g^2 pN\phi^2/f_{\text{T}}^{p+1} \quad 2g^2 pN\phi^2 f_p/f_{\text{T}}^{p+2}$

$$\kappa_+ = \kappa = pNf_p/f_{\text{T}}^{p+2} \quad \text{and} \quad \kappa_- = pN/f_{\text{T}}^{p+1}, \quad (17)$$

from which κ_+ equals the double eigenvalue κ of $M_{\Phi\Phi}$ for $K = \tilde{K}_{2\text{st}}$. Inserting Eqs. (12) and (15) in the second term of the right-hand side of Eq. (9) we can define the canonically normalized fields that are denoted by a hat. Combining Eqs. (16) and (17) we can establish for both $\tilde{K}_{\ell\text{st}}$ in Eq. (4) a unique expression regarding the normalized inflaton $\hat{\phi}$, i.e.,

$$d\hat{\phi}/d\phi = J = \sqrt{2\kappa} \quad \text{for } \phi < 1. \quad (18)$$

This result is identical to the $\phi - \hat{\phi}$ relation obtained for $K = K_{1\text{T}}$ (or $\tilde{K}_{1\text{T}}$) in Ref. [47]. For the remaining fields of the $\bar{\Phi} - \Phi$ system we find

$$\begin{aligned} \hat{\theta}_\pm &= \sqrt{\kappa_\pm} \phi \theta_\pm, & \hat{\theta}_\Phi &= \sqrt{2\kappa_-} \phi (\theta_\Phi - \pi/4) \quad \text{for } K = \tilde{K}_{1\text{st}}, \\ \hat{\theta}_\pm &= \sqrt{\kappa} \phi \theta_\pm, & \hat{\theta}_\Phi &= \sqrt{2\kappa} \phi (\theta_\Phi - \pi/4) \quad \text{for } K = \tilde{K}_{2\text{st}}, \end{aligned}$$

where $\theta_\pm = (\bar{\theta} \pm \theta)/\sqrt{2}$ —needless to say, $(s, \bar{s})$ are canonically normalized since $\langle K_{SS^*} \rangle_I = 1$. Note, in passing, that the spinors ψ_S and $\psi_{\Phi\pm}$ associated with the superfields S and $\Phi - \bar{\Phi}$ are normalized similarly, i.e., $\hat{\psi}_{\Phi+} = \sqrt{\kappa} \psi_{\Phi+}$ with $\psi_{\Phi\pm} = (\psi_\Phi \pm \psi_{\bar{\Phi}})/\sqrt{2}$.

We can verify that the inflationary direction in Eq. (13) is stable with respect to the fluctuations of the fields $\chi^\alpha = \theta_+, \theta_\Phi$ and S . To this end, we construct the mass-squared spectrum of those scalars. Taking the limit $M \ll 1$ we find the expressions of the masses squared $m_{\chi^\alpha}^2$ arranged in Table I, which approach rather well the exact expressions taken into account in our numerical computation. These expressions assist us to appreciate the role of $N_{\text{st}} > 0$ in retaining positive m_s^2 . Moreover, $m_{\chi^\alpha}^2 \gg H_{\text{HI}}^2 = V_{10}/3$ for $\phi_{\text{f}} \leq \phi \leq \phi_\star$ —where $\phi_\star$ and ϕ_{f} are the values of ϕ when $k_\star = 0.05/\text{Mpc}$ crosses the horizon of $p\text{HI}$ and at the end of $p\text{HI}$ correspondingly. In Table I we display also the masses M_{BL} of the gauge boson A_{BL} and the masses of the corresponding fermions with the eigenstate $\hat{\psi}_\pm$ defined as

$\hat{\psi}_{\pm} = (\hat{\psi}_{\Phi+} \pm \psi_S)/\sqrt{2}$. Since $\mathbb{G}_{B-L}$ is broken during p HI, no cosmic strings are formed after its termination.

The derived mass spectrum can be employed in order to find the one-loop radiative corrections ΔV_{HI} to V_{HI} . Considering SUGRA as an effective theory with cutoff scale equal to m_p , the well-known Coleman-Weinberg formula can be employed self-consistently taking into account only the masses that lie well below m_p , i.e., all the masses arranged in Table I besides M_{BL} and m_{θ_ϕ} —cf. Refs. [62,63]. The resulting ΔV_{HI} keeps intact our inflationary outputs, provided that the renormalization-group mass scale Λ is determined by requiring $\Delta V_{\text{HI}}(\phi_*)=0$ or $\Delta V_{\text{HI}}(\phi_f)=0$; e.g., imposing the former condition for $N_{\text{st}}=1$, we find $\Lambda \simeq (2.2\text{--}5.1) \times 10^{-5}$ for the allowed parameter space exposed below. Under these circumstances, our inflationary predictions can be exclusively reproduced by using V_{HI} in Eq. (14).

IV. INFLATIONARY REQUIREMENTS

Applying the standard formulas quoted in Ref. [47] for $V_I = V_{\text{HI}}$, we can compute a number of observational quantities, which assist us to qualify our inflationary setting. Namely, we extract the number N_* of e-foldings that the scale k_* experiences during p HI and the amplitude A_s of the power spectrum of the curvature perturbations generated by ϕ for $\phi = \phi_*$. These observables must be compatible with the requirements [2]

$$N_* \simeq 61.3 + \frac{1-3w_{\text{rh}}}{12(1+w_{\text{rh}})} \ln \frac{\pi^2 g_{\text{rh}*} T_{\text{rh}}^4}{30 V_{\text{HI}}(\phi_f)} + \frac{1}{4} \ln \frac{V_{\text{HI}}(\phi_*)^2}{g_{\text{rh}*}^{1/3} V_{\text{HI}}(\phi_f)} \quad \text{and} \quad A_s^{\frac{1}{2}} \simeq 4.618 \times 10^{-5}, \quad (19)$$

where we take into account that p HI is followed, in turn, by an oscillatory phase with mean equation-of-state parameter w_{rh} , radiation, and matter domination. We estimate w_{rh} as detailed in the Appendix and we find $w_{\text{rh}} \simeq (0.12\text{--}0.3)$ in the major part of the parameter space—note that w_{rh} decreases with N .

The constraints above allow us to extract λ and ϕ_* for any given p , N , and M . However, M is irrelevant for the inflationary observables, provided that it is confined to values much lower than m_p . Motivated by the unification of the gauge coupling constants with the particle content of the MSSM, we determine M identifying the unification scale $M_{\text{GUT}} \simeq 2/2.433 \times 10^{-2}$ with the value of M_{BL} —see Table I—at the SUSY vacuum in Eq. (3). Given that $\langle f_T \rangle \simeq 1$ for $M \ll 1$, we obtain, for both $\tilde{K}_{\ell\text{st}}$ in Eq. (4),

$$M \simeq M_{\text{GUT}}/g\sqrt{2Np} \quad (20)$$

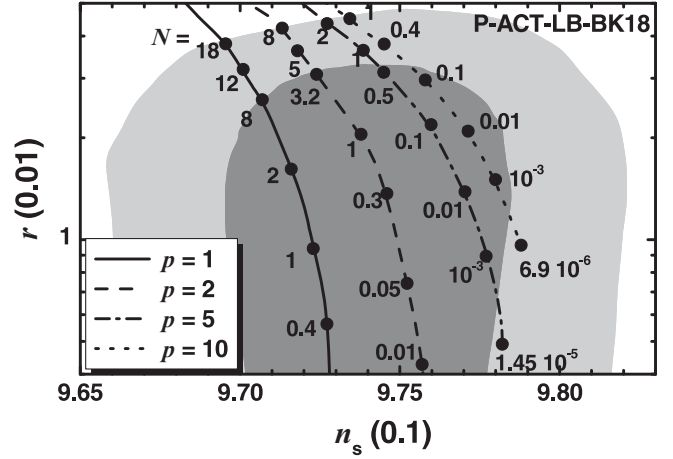


FIG. 1. Allowed curves in the n_s - r plane for various p 's shown in the legend and varying N as shown along the curves. The marginalized joint 68% (95%) regions from P-ACT-LB-BK18 data are depicted by the dark (light) shaded contours.

with $g \simeq 0.7$ being the value of the GUT gauge coupling constant. The success of our inflationary setting is assured if

$$M \lesssim m_p \Rightarrow N \gtrsim N_{\min} = M_{\text{GUT}}^2/2g^2 p m_p^2 \quad (21)$$

and so a lower bound on N is applied. We are then left with only two free input parameters (p, N), which can be constrained computing n_s , its running a_s , and r . These quantities can be found approximately by applying the analytic expressions displayed in Ref. [47]. Namely, we found

$$n_s \simeq 1 - \frac{p+2}{(p+1)N_*}, \quad a_s \simeq -\frac{p+2}{(p+1)N_*^2}, \quad (22a)$$

which are independent of the exponent $n=4$ of ϕ in Eq. (14), whereas r for the specific n value is written as

$$r \simeq 4(4^p p N)^{1/(p+1)}/(p+1)N_*^{(p+2)/(p+1)}. \quad (22b)$$

The expressions above for n_s and a_s are independent from N in the lowest order expansion in powers of $1/N_*$. This feature, however, is not totally kept, as we show numerically in Fig. 1, where we compare our outputs against the P-ACT-LB-BK18 data [2] in the n_s - r plane. We draw solid, dashed, dot-dashed, and dotted lines for $p=1, 2, 5$, and 10 , respectively, and show the variation of N along each line. For the two latter p 's, the corresponding $N_{\min}$'s from Eq. (21) are also displayed. We clearly see that n_s increases with p and r with N and so the whole shaded region favored by P-ACT-LB-BK18 is covered varying p and N . As a bonus, the resulting r 's for (p, N) 's of order

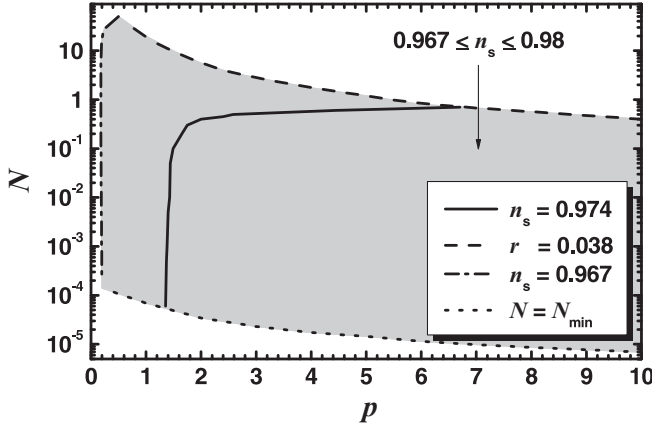


FIG. 2. Allowed (shaded) region as determined by Eqs. (1), (19), and (21) in the p - N plane. Along the solid line, we fix $n_s = 0.974$. The conventions adopted for the various lines are also shown.

unity lie within the reach of planned experiments [66] aiming to discover primordial gravitational waves.

The variation of the observables in Fig. 1 reveals that our free parameters can be constrained in the p - N plane. The relevant allowed region is shown in Fig. 2. It is bounded by (i) the dashed line, which originates the bound on r in Eq. (1), (ii) the dot-dashed line, which comes from the lower bound on n_s in Eq. (1), and (iii) the dotted line along which Eq. (21) is saturated—we take by hand a maximal $p = 10$ since no upper bound on p can be inferred from the corresponding bound on n_s in Eq. (1). Fixing n_s to its central value in Eq. (1), we obtain the solid line along which we obtain

$$1.355 \lesssim p \lesssim 6.7 \quad \text{and} \quad 6 \times 10^{-5} \lesssim N \lesssim 0.7 \quad (23)$$

with $N_\star \simeq (54.7-56.8)$, $a_s \simeq -4.7 \times 10^{-4}$, and $r \gtrsim 2.8 \times 10^{-4}$. The obtained $|a_s|$'s remain negligibly small being, thereby, consistent with its 95% CL allowed margin in Ref. [2].

Note, finally, that our proposal is stabilized against corrections from higher order terms—e.g., $(\bar{\Phi}\Phi)^l$ with $l > 1$ in Eq. (2)—since $\phi_\star < 1$, as dictated by Eq. (18). The closer to unity $\phi_\star$ is chosen, the largest $N_\star$ is obtained. To quantify the relevant tuning, we estimate $\Delta_\star = 1 - \phi_\star$ which ranges in the interval (0.8–28)% for the values in Eq. (23). The maximal $\Delta_\star$ values are obtained for the largest p 's and N 's indicating, thereby, that naturalness prefers N and p values of order unity.

V. POSTINFLATIONARY COMPLETION

Our inflationary setting can be easily embedded in a $B-L$ extension of the MSSM promoting to gauge the pre-existing global $U(1)_{B-L}$ —cf. Refs. [53,63]. The terms of the total super- and Kähler potentials that control the

TABLE II. Mass-squared spectrum of the noninflaton sector along the path in Eqs. (13) and (25) for $K = \tilde{K}_{\text{st}} + \Delta K$ with $\ell = 1, 2$.

Fields	Eigenstates	Masses squared
22 real scalar	$h_\pm, \bar{h}_\pm$	$m_{h_\pm}^2$
	$\tilde{\nu}_i^c, \bar{\tilde{\nu}}_i^c$	$m_{\tilde{\nu}^c}^2$
	$\tilde{l}_i, \bar{\tilde{l}}_i$	$m_{\tilde{l}}^2$
Three Weyl spinors	N_i^c	$M_{iN^c}^2$
		$3H_{\text{HI}}^2(1 + 1/N_{\text{st}}) \pm 2\lambda_\mu/\lambda\phi^2$
		$3H_{\text{HI}}^2(1 + 1/N_{\text{st}} + 8\lambda_{iN^c}^2/\lambda^2\phi^2)$
		$3H_{\text{HI}}^2(1 + 1/N_{\text{st}})$
		$24\lambda_{iN^c}^2 H_{\text{HI}}^2/\lambda^2\phi^2$

coexistence of the inflationary and the MSSM sectors are

$$\Delta W = \lambda_\mu S H_u H_d + h_{ijN} N_i^c L_j H_u + \lambda_{iN^c} \bar{\Phi} N_i^{c2}; \quad (24a)$$

$$\Delta K = N_{\text{st}} \ln \left(1 + \sum_{a=1}^4 |X_a|^2 / N_{\text{st}} \right) \quad (24b)$$

with $X_\alpha = \tilde{N}_i^c, \tilde{L}_j, H_u$, and H_d , where $\tilde{N}_i^c$ and $\tilde{L}_j$ are the scalar superpartners of the right-handed neutrinos N_i^c and left-handed leptons L_j , respectively. Here we adopt the notation and the $B-L$ and R charges of the various superfields as displayed in Table 1 of Ref. [63]. For simplicity, we use the same prefactors of $\ln$ in all terms of ΔK .

We assume that X_α are stabilized at zero during p HI and so the inflationary trajectory in Eq. (13) has to be supplemented by the condition

$$\langle H_u \rangle_I = \langle H_d \rangle_I = \langle \tilde{N}_i^c \rangle_I = \langle \tilde{L}_j \rangle_I = 0. \quad (25)$$

The stability of this path can be checked, parametrizing the complex fields above as we do for S in Eq. (12)—note that we use small letters for the real and imaginary components. The relevant masses squared are listed in Table II where we see that $m_{\tilde{\nu}^c}^2 > 0$, $m_{\tilde{l}}^2 > 0$ and $m_{h_\pm}^2 > 0$ for any $\phi_f < \phi < 1$ —here $m_{h_\pm}$ are the eigenvalues associated with the Higgs eigenstates

$$h_\pm = (h_u \pm h_d)/\sqrt{2} \quad \text{and} \quad \bar{h}_\pm = (\bar{h}_u \pm \bar{h}_d)/\sqrt{2}. \quad (26)$$

On the other hand, the positivity of $m_{h_-}^2$ dictates the establishment of the inequality (cf. Ref. [53]),

$$\lambda_\mu < \lambda(1 + N_{\text{st}})\phi_f^2/2N_{\text{st}}. \quad (27)$$

We can verify numerically that the requirement above is fulfilled for $N_{\text{st}} = 1$ and $\lambda_\mu \lesssim 2 \times 10^{-5}$. Despite this low upper bound, λ_μ can become compatible with $\mu/m_{3/2} \sim 1$ as we show in Sec. VA, where the generation of the μ -term of the MSSM is analyzed. On the other hand, the implementation of nTL is explored in Sec. VB and our final results, including the postinflationary constraints,

are given in Sec. V C. Note in passing that, due to the extra contributions of Table II, ΔV_{HI} increases and so Λ should be readjusted. For example, for $p = 2$ we find $\Lambda \simeq (4-5) \times 10^{15}$ GeV—hereafter we restore units, i.e., we take $m_{\text{P}} = 2.43 \times 10^{18}$ GeV.

A. Generation of the μ -term of the MSSM

The origin of the μ -term of the MSSM can be explained if we combine the terms of Eq. (2) with the first term in Eq. (24a) working at the level of SUSY. In fact, the total low-energy potential is

$$V_{\text{tot}} = V_{\text{F}} + V_{\text{soft}}, \quad (28)$$

where the first term includes the SUSY F-term contributions into V_{tot} which read

$$\begin{aligned} V_{\text{F}} &= K_{\text{SUSY}}^{\alpha\bar{\beta}} (W + \Delta W)_{\alpha} (W + \Delta W)_{\bar{\beta}}^* \\ &= \lambda^2 |\bar{\Phi}\Phi - M^2/2|^2 + \lambda^2 |S|^2 (|\bar{\Phi}|^2 + |\Phi|^2)/Np + \dots \end{aligned} \quad (29)$$

where K_{SUSY} can be obtained by expanding $\tilde{K}_{\ell\text{st}} + \Delta K$ —see Eqs. (4) and (24b)—for low field values with result

$$K_{\text{SUSY}} = Np(|\Phi|^2 + |\bar{\Phi}|^2 - \bar{\Phi}\Phi - \bar{\Phi}^*\Phi^*) + |S|^2 + |X^\alpha|^2. \quad (30)$$

Also, the ellipsis in Eq. (29) includes terms that vanish at the SUSY vacuum due to the zero VEVs of H_u , H_d , and $\tilde{N}_i^c$. On the other hand, the contribution into V_{tot} from the soft SUSY-breaking terms reads

$$V_{\text{soft}} = (\lambda A_\lambda S \bar{\Phi}\Phi - \lambda a_S M^2 S/2 + \text{H.c.}) + m_\gamma^2 |X^\gamma|^2, \quad (31)$$

where $X^\gamma = S, \Phi, \bar{\Phi}, H_u$, and H_d . Also m_γ , A_λ , and a_S are soft SUSY-breaking mass parameters much lower than M . Rotating S in the real axis by an appropriate R transformation, choosing conveniently the phases of A_λ and a_S —so as V_{tot} in Eq. (28) to be minimized—and substituting the Φ and $\bar{\Phi}$ values from Eq. (3), we get

$$\langle V_{\text{tot}}(S) \rangle = \lambda^2 M^2 S^2 / Np - 2\lambda a_\mu m_{3/2} M^2 S, \quad (32a)$$

where $m_{3/2}$ is the $\tilde{G}$ (gravitino) mass and

$$a_\mu = (|A_\lambda| + |a_S|)/2m_{3/2} > 0 \quad (32b)$$

is a parameter of order unity that parametrizes our ignorance about the dependence of $|A_\lambda|$ and $|a_S|$ on $m_{3/2}$. The extermination condition of $\langle V_{\text{tot}}(S) \rangle$ with respect to S shifts slightly $\langle S \rangle$ from zero in Eq. (3)—cf. Ref. [67]—as follows:

$$d\langle V_{\text{tot}}(S) \rangle / dS = 0 \Rightarrow \langle S \rangle = 2Np a_\mu m_{3/2} / \lambda. \quad (32c)$$

The generated μ parameter from the first term in Eq. (24a),

$$\mu = \lambda_\mu \langle S \rangle = 2\lambda_\mu Np a_\mu m_{3/2} / \lambda, \quad (33)$$

can be comparable to $m_{3/2}$ for $N \sim 1$, if λ_μ is of the same order of magnitude with λ . Since $\lambda \sim 10^{-6}$, due to the $\lambda - A_s$ condition implied by the rightmost expression in Eq. (19)—see Ref. [47]—the resulting λ_μ can become compatible with Eq. (27) assuring, thereby, the stability of the path in Eq. (25) during $p\text{HI}$.

B. Nonthermal leptogenesis and $\tilde{G}$ constraint

In addition to the generation of the μ -term, ΔW in Eq. (24a) allows for the implementation of nTL via the direct decay of the inflaton into N_i^c . For the success of this process, however, one has to restrict the decay channel of the inflaton into H_u and H_d which opens via the first term of ΔW . This implies constraints on λ_μ or, via Eq. (33), on the ratio $\mu/m_{3/2}$.

More explicitly, when $p\text{HI}$ is over, the (canonically normalized) inflaton

$$\widehat{\delta\phi} = \langle J \rangle \delta\phi \quad \text{with} \quad \delta\phi = \phi - M \quad \text{and} \quad \langle J \rangle \simeq \sqrt{2Np} \quad (34)$$

continues to roll down toward the SUSY vacuum, Eq. (3), where it acquires mass—reducing with increasing N and p —given by

$$\hat{m}_{\delta\phi} = \langle V_{\text{HI}, \hat{\phi}} \rangle^{1/2} \simeq \lambda M / \sqrt{Np}. \quad (35)$$

After a brief stage of tachyonic preheating [75] that does not lead to significant particle production [76], $\widehat{\delta\phi}$ settles into a phase of damped oscillations initially around zero—where V_{HI} has a maximum—and then around one of the minima of V_{HI} similar to the picture described in Appendix B of Ref. [59] for the nonminimal HI. Whenever $\widehat{\delta\phi}$ passes through zero, several particles (mostly A_{BL} bosons) may be produced via the mechanism of instant preheating [77]. This process becomes more efficient as N decreases since the amplitude of the first oscillation

$$\phi_{\text{mx}} \simeq \sqrt{6/(16 + 3Np)} m_{\text{P}} \quad (36)$$

increases, increasing, thereby, the frequency of the violation of the adiabaticity criterion [76,77]. Given that the efficiency of this mechanism requires a rather large (let say 50–70) number of passages of $\widehat{\delta\phi}$ through zero, we intuitively expect that our estimates below are more accurate for $N > 0.01$, which ensures a less frequent passage of $\widehat{\delta\phi}$ through zero weakening, thereby, the effects from preheating.

Nonetheless, the standard perturbative approach to the $\widehat{\delta\phi}$ decay provides a very efficient total decay rate

$$\hat{\Gamma}_{\delta\phi} = \hat{\Gamma}_{\delta\phi \rightarrow N_i^c} + \hat{\Gamma}_{\delta\phi \rightarrow H}, \quad (37)$$

where the individual decay widths are

$$\hat{\Gamma}_{\delta\phi \rightarrow N_i^c} = (g_{iN^c}^2 \hat{m}_{\delta\phi} / 16\pi) (1 - 4M_{iN^c}^2 / \hat{m}_{\delta\phi}^2)^{3/2}; \quad (38a)$$

$$\hat{\Gamma}_{\delta\phi \rightarrow H} = 2g_H^2 \hat{m}_{\delta\phi} / 8\pi \quad (38b)$$

with the relevant coupling constants being

$$g_{iN^c} = \sqrt{2}\lambda_{iN^c} / \langle J \rangle \quad \text{and} \quad g_H = \lambda_\mu / \sqrt{2}. \quad (38c)$$

Moreover, M_{iN^c} are the Majorana masses of N_i^c identified by the third term in Eq. (24a) as

$$M_{iN^c} = \sqrt{2}\lambda_{iN^c} M \quad \text{with} \quad \lambda_{iN^c} \leq 4\pi \quad (39)$$

to be consistent with the perturbative bound. The Lagrangian terms that describe the decay channels above are similar to those given, e.g., in Ref. [63]. Note that three-particle decay of $\widehat{\delta\phi}$ is suppressed as noted there and possible perturbative decay into A_{BL} bosons is kinetically blocked. Thanks to $\hat{\Gamma}_{\delta\phi}$, the Universe is reheated at a temperature [78,79]

$$T_{\text{rh}} = \left(\frac{90}{\pi^2 g_{\text{rh}*}} \right)^{1/4} \left(\frac{2(w_{\text{rh}} + 1)}{5 - 3w_{\text{rh}}} m_p \hat{\Gamma}_{\delta\phi} \right)^{1/2} \quad (40)$$

with $g_{\text{rh}*} = 228.75$ the MSSM effective number of relativistic degrees of freedom.

For $M_{iN^c} T_{\text{rh}}$, the out-of-equilibrium decay of N_i^c via the second term in Eq. (24a) generates a L -asymmetry yield that is partially converted through sphaleron effects into a yield of the observed B asymmetry of the Universe [68,70],

$$Y_B = -0.35 \cdot 2 \cdot \frac{5 - 3w_{\text{rh}}}{4(1 - w_{\text{rh}})} \frac{T_{\text{rh}}}{\hat{m}_{\delta\phi}} \frac{\hat{\Gamma}_{\delta\phi \rightarrow N_i^c}}{\hat{\Gamma}_{\delta\phi}} \varepsilon_L. \quad (41)$$

The simplified formula above can be derived assuming normal hierarchy for M_{iN^c} (i.e., $M_{1N^c} \ll M_{2N^c}, M_{3N^c}$) and that $\widehat{\delta\phi}$ decays via $\hat{\Gamma}_{\delta\phi \rightarrow N_i^c}$ in Eq. (38a) exclusively into N_1^c . Under these circumstances, we can obtain a maximal value for the L asymmetry ε_L which is [69,80]

$$\varepsilon_L = -\frac{3}{8\pi} \frac{m_{\nu_\tau} M_{1N^c}}{\langle H_u \rangle^2}, \quad \text{where} \quad m_{\nu_\tau} = \sqrt{\Delta m_{\oplus}^2} \simeq 0.05 \text{ eV} \quad (42)$$

is the mass of heaviest light neutrino ν_τ , and $\Delta m_{\oplus}^2$ [81] is the atmospheric neutrino mass-squared difference. Also,

we set $\langle H_u \rangle = 174 \text{ GeV}$, adopting the large $\tan\beta$ regime of the MSSM. The validity of Eq. (41) requires the kinematic allowance of the decay $\widehat{\delta\phi} \rightarrow 2N_1^c$ and the avoidance of any erasure of the produced lepton-asymmetry yield. These conditions can be fulfilled if we respectively demand

$$\hat{m}_{\delta\phi} \geq 2M_{1N^c} \quad \text{and} \quad M_{1N^c} \gtrsim 10T_{\text{rh}}. \quad (43)$$

The result of Eq. (41) has to meet the observational one [1]

$$Y_B = (8.75 \pm 0.085) \times 10^{-11} \quad \text{at 95\% CL.} \quad (44)$$

On the other hand, T_{rh} must be compatible with constraints on the $\tilde{G}$ abundance, $Y_{3/2}$, at the onset of big bang nucleosynthesis (BBN), which is estimated to be [72]

$$Y_{3/2} \simeq 2.9 \times 10^{-13} (T_{\text{rh}} / \text{EeV}) / (3 - w_{\text{rh}}), \quad (45)$$

where we take into account only thermal production of $\tilde{G}$ and assume that $\tilde{G}$ is much heavier than the MSSM gauginos. On the other hand, $Y_{3/2}$ is bounded from above in order to avoid spoiling the success of the BBN. For the typical case where $\tilde{G}$ decays with a tiny hadronic branching ratio, the bounds on $Y_{3/2}$ imply the following upper bounds on T_{rh} [71,72]:

$$T_{\text{rh}} \lesssim \begin{cases} 0.53 \text{ EeV} \\ 5.3 \text{ EeV} \end{cases} \quad \text{for } m_{3/2} \simeq \begin{cases} 10.6 \text{ TeV}, \\ 13.6 \text{ TeV}, \end{cases} \quad (46)$$

at 95% CL. The bounds above can be somehow relaxed in the case of a stable $\tilde{G}$ (cf. Ref. [41]).

C. Results

The compatibility of the restrictions in Eqs. (27), (43), (44), and (46) together with those from Eqs. (1) and (19) is illustrated in Fig. 3 where we present allowed curves in N - M_{1N^c} and N - T_{rh} planes for $p = 2$, $N_{\text{st}} = 1/3$, and $\mu/m_{3/2} = 0.01, 0.1, 1$, and 3 . The chosen p value is very close to the p values that yield the central n_s in Eq. (1), as shown in Fig. 2, and the one for N_{st} assists us to overcome the constraint of Eq. (27) without extensive tuning. Moreover, the considered $\mu/m_{3/2}$ ratios are encountered in various recent studies [82–84] of the parameter space of the MSSM with $m_{3/2}$ values comparable to those considered in Eq. (46). Despite the fact that the obtained μ value is quite low for $\mu/m_{3/2} = 0.01$, it is, in general, still allowed by the relevant experimental exclusion limit [85] on the chargino mass.

The ratio $\mu/m_{3/2}$ controls the λ_μ value derived by Eq. (33) and via it the contribution of $\hat{\Gamma}_{\delta\phi \rightarrow H}$ in Eq. (38b) into the $\hat{\Gamma}_{\delta\phi}$ —see Eq. (37). As expected, the M_{1N^c} required by Eq. (44) increases with $\mu/m_{3/2}$ in order to reinforce the contribution of $\hat{\Gamma}_{\delta\phi \rightarrow N_i^c}$ [see Eq. (38a)] into

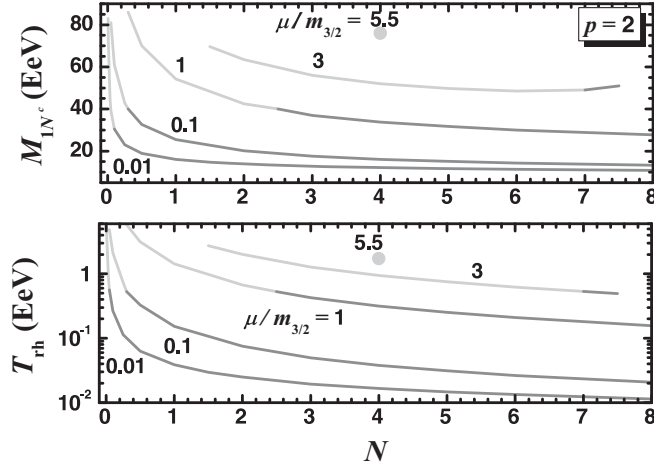


FIG. 3. Contours yielding the central Y_B in Eq. (44) consistent with Eqs. (1), (19), (27), (43), and (46) in the N - M_{1N^c} and N - T_{rh} planes for $p = 2$, $N_{st} = 1/3$, and various $\mu/m_{3/2}$ values indicated on the curves—recall that 1 EeV = 10^9 GeV. The gray (light gray) segments of the lines are allowed by Eq. (46) for $m_{3/2} = 10.6$ TeV ($m_{3/2} = 13.6$ TeV).

$\hat{\Gamma}_{\delta\phi}$. The increase of $\mu/m_{3/2}$ is constrained directly by Eq. (27) and indirectly by Eq. (46) since it causes an increase of T_{rh} [see Eq. (40)]. The gray (light gray) segments of the lines displayed in Fig. 3 are allowed by Eq. (46) for $m_{3/2} = 10.6$ TeV ($m_{3/2} = 13.6$ TeV). As shown from the lower plot in Fig. 3, the lines corresponding to $\mu/m_{3/2} = 0.01$ and 0.1 terminate for low N values due to the upper bound on T_{rh} in Eq. (46). On the other hand, for $\mu/m_{3/2} = 1$ and 3 this termination occurs due to the saturation of the bound on λ_μ in Eq. (27). The $\mu/m_{3/2} = 3$ line is cut for $N = 7.5$ due to the leftmost inequality in Eq. (43). The competition of these two aforementioned constraints causes a progressive reduction of the length of the allowed contours as $\mu/m_{3/2}$ increases beyond 3 until they shrink to one point for $\mu/m_{3/2} \simeq 5.5$, corresponding to $N = 4$, $M_{1N^c} = 76$ EeV, and $T_{rh} = 1.73$ EeV. Along the curves of Fig. 3, we obtain

$$0.035 \lesssim \lambda_\mu/10^{-6} \lesssim 40 \quad \text{and} \quad 0.14 \lesssim \hat{m}_{\delta\phi}/\text{ZeV} \lesssim 32, \quad (47)$$

where 1 ZeV = 10^{12} GeV. The largest λ_μ is obtained for $\mu/m_{3/2} = 1$ and $N = 0.3$, whereas the lowest one is found for $\mu/m_{3/2} = 0.01$ and $N = 8$. On the other hand, $\hat{m}_{\delta\phi}$ decreases with increasing N independent from $\mu/m_{3/2}$. Had we used larger N_{st} values, most of the curves with $\mu/m_{3/2} > 1$ would have been excluded. Therefore, our scheme primarily prefers the hierarchy $\mu \leq m_{3/2}$, which is defended by various analyses [82–84] of the MSSM parameter space.

As a bottom line, nTL is a realistic possibility within our setting since it can be reconciled with the $\tilde{G}$ constraint for $m_{3/2} \sim 10$ TeV and several MSSM versions.

VI. CONCLUSIONS

Prompted by the P-ACT-LB-BK18 data, we updated pole [52] (or T-model [53]) Higgs inflation by extending our proposal in Ref. [47]. Namely, we adopted the super- and Kähler potentials W and K in Eqs. (2) and (4) which lead to the inflationary potential in Eq. (14) and the kinetic mixing in Eq. (18). Moreover, we required that the Higgs inflaton assumes ultimately its VEV predicted by the gauge unification within the MSSM. This requirement influenced via Eq. (21) the parameter space of pHI and via Eqs. (35) and (39) the postinflationary scenario. Confining (p, N) in the ranges shown in Fig. 2 we achieved a nice covering of the present data—see Fig. 1. Fixing n_s to its central value in Eq. (1) we obtained the clear predictions of Eq. (23) regarding our free parameters. As a consequence, our present framework drastically reduces the allowed (p, N) domain found for a gauge-singlet inflaton in Ref. [47]. It is gratifying that the resulting r values for (p, N) 's of order unity could be observationally accessible in the foreseeable future. We also specified a postinflationary completion, based on the extra W - and K -terms in Eqs. (24a) and (24b), which offers a nice solution to the μ problem of the MSSM and allows for baryogenesis via nTL with M_{1N^c} in the range (10–76) EeV provided that $\mu/m_{3/2} \leq 5.5$.

Since our main aim here was the introduction of pHI, we opted to utilize its embedding in a GUT based on $\mathbb{G}_{B-L}$. Obviously our scheme can be realized, employing the same W and K 's, within other SUSY GUTs too, based on more structured gauge groups provided that Φ and $\bar{\Phi}$ consist of a conjugate pair of Higgs superfields—cf. Refs. [59,61]. Special treatment is required if the Higgs superfields belong to the adjoint representation of a GUT—cf. Ref. [65]. Moreover, given that the term $\lambda M^2 S/4$ of W in Eq. (2) plays no role during pHI, our scenario can be implemented by replacing it with $\lambda' S^3$ and identifying Φ and $\bar{\Phi}$ with the electroweak Higgs doublets H_u and H_d of the next to MSSM [74]. In the last case, a tighter connection of the high- to the low-energy phenomenology is expected.

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX: REHEATING PROCESS, LEPTON ASYMMETRY, AND GRAVITINO ABUNDANCES

We present here a numerical, in Appendix A 1, and analytic, in Appendix A 2, description of the postinflationary evolution of our setting, justifying, thereby, the

expressions adopted in Eqs. (40), (41), and (45). From them, the two latter results are presented for first time, to our knowledge, in the literature.

1. Formulation of the postinflationary dynamics

We display here the equations that govern the evolution of the various energy and number densities involved in our scenario of nTL. The initial form of these equations is given in Appendix A 1 a, whereas a form more convenient for numerical manipulations is given in Appendix A 1 b.

a. Initial form

The energy densities ρ of $\widehat{\delta\phi}$, the energy density ρ_R of the produced radiation, and the number densities n_L of the lepton asymmetry and $n_{3/2}$ of the $\tilde{G}$'s satisfy the following Boltzmann equations—cf. Ref. [86]:

$$\dot{\rho} + 3(w_{\text{rh}} + 1)H\rho + (w_{\text{rh}} + 1)\hat{\Gamma}_{\delta\phi}\rho = 0, \quad (\text{A1a})$$

$$\dot{\rho}_R + 4H\rho_R - (w_{\text{rh}} + 1)\hat{\Gamma}_{\delta\phi}\rho = 0, \quad (\text{A1b})$$

$$\dot{n}_L + 3Hn_L - 2\varepsilon_L(w_{\text{rh}} + 1)\hat{\Gamma}_{\delta\phi \rightarrow N_i^c}\rho/\hat{m}_{\delta\phi} = 0, \quad (\text{A1c})$$

$$\dot{n}_{3/2} + 3Hn_{3/2} - C_{3/2}n_{\text{eq}}^2 = 0. \quad (\text{A1d})$$

Here the overdot denotes derivation with respect to the cosmic time t . Because of the polynomial character of V_{HI} in Eq. (14) and the nonminimal kinetic mixing in Eq. (18), the estimation of w_{rh} requires some care—cf. Ref. [87]. We determine it adapting the general formula [86], i.e.,

$$w_{\text{rh}} = 2 \frac{\int_M^{\phi_{\text{mx}}} d\phi J(1 - V_{\text{HI}}/V_{\text{mx}})^{1/2}}{\int_M^{\phi_{\text{mx}}} d\phi J(1 - V_{\text{HI}}/V_{\text{mx}})^{-1/2}} - 1, \quad (\text{A2})$$

where $V_{\text{mx}} = V_{\text{HI}}(\phi_{\text{mx}})$. The amplitude of the oscillations during reheating ϕ_{mx} is found by solving numerically the condition $\sqrt{3}H_{\text{HI}}(\phi_{\text{mx}}) = \hat{m}_{\delta\phi}$ if $\sqrt{3}H_{\text{HI}}(\phi_f) > \hat{m}_{\delta\phi}$ or it is $\phi_{\text{mx}} = \phi_f$ otherwise. For $N = (0.1-8)$ where $J \simeq 1$, w_{rh} can be well approximated by performing an expansion for low M values with the following result:

$$w_{\text{rh}} \simeq \frac{1}{3} - 0.51 \frac{M}{\phi_{\text{mx}}} + 0.22 \frac{M^2}{\phi_{\text{mx}}^2} + \dots, \quad (\text{A3})$$

which implies a reduction of w_{rh} with respect to its value (1/3) corresponding to a pure quartic potential [86]. We also checked that $V_{\text{mx}} \gg V_{\text{HI}}(\phi = 0)$ and so the oscillations about the central maximum of V_{HI} (for $\phi = 0$) cannot be ignored.

Moreover, n_{eq} in Eq. (A1d) is the equilibrium number density of each bosonic relativistic species

$$n_{\text{eq}} = \zeta(3)T^3/\pi^2 \quad (\text{A4})$$

and $C_{3/2}$ is a collision term for $\tilde{G}$ production which, in the limit of massless MSSM gauginos, turns out to be [71,72]

$$C_{3/2} = \frac{3\pi}{16\zeta(3)m_{\text{P}}^2} \sum_{\alpha=1}^3 c_{\alpha} g_{\alpha}^2 \ln\left(\frac{k_{\alpha}}{g_{\alpha}}\right), \quad (\text{A5})$$

where $(c_{\alpha}) = (33/5, 27, 72)$, g_{α} are the gauge coupling constants of the MSSM [$(g_{\alpha}(T_{\text{rh}})) = (0.54, 0.68, 0.88)$ for $T_{\text{rh}} \sim \text{EeV}$], and $(k_{\alpha}) = (1.634, 1.312, 1.271)$. Clearly, in the limit of massless MSSM gauginos, the resulting $n_{3/2}$ is practically $m_{3/2}$ independent.

Finally, the Hubble expansion parameter H during this period is given by

$$H = (\rho + \rho_R)^{1/2}/\sqrt{3}m_{\text{P}}, \quad (\text{A6})$$

whereas the temperature T and the entropy density s are found from the relations

$$\rho_R = \frac{\pi^2}{30}g_*T^4 \quad \text{and} \quad s = \frac{2\pi^2}{45}g_*T^3. \quad (\text{A7})$$

The system of Eqs. (A1a)–(A1d) is solved under the following initial conditions:

$$\rho(0) = V_{\text{HI}}(\phi_{\text{mx}}) \quad \text{and} \quad \rho_R(0) = n_L(0) = n_{3/2}(0) = 0. \quad (\text{A8})$$

The results are obtained for a temperature $T_f \ll T_{\text{rh}}$, where T_{rh} is defined from the condition [78,79]

$$\rho(T_{\text{rh}}) = \rho_R(T_{\text{rh}}). \quad (\text{A9})$$

The abundances in Eqs. (41) and (45) are estimated by the solution of Eqs. (A1c) and (A1d) via the expressions

$$Y_L = (n_L/s)(T_f) \quad \text{and} \quad Y_{3/2} = (n_{3/2}/s)(T_f). \quad (\text{A10})$$

b. Reformulation

The numerical integration of Eqs. (A1a)–(A1d) is facilitated, if we define the following variables—cf. Ref. [78]:

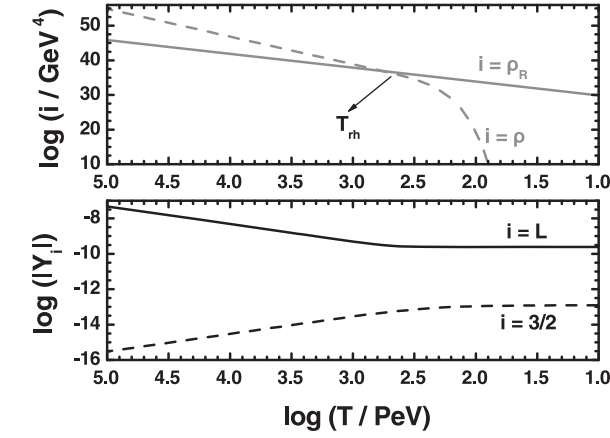
$$f = \rho R^{3(1+w_{\text{rh}})}, \quad f_R = \rho_R R^4, \quad (\text{A11a})$$

$$f_L = n_L R^3 \quad \text{and} \quad f_{3/2} = n_{3/2} R^3, \quad (\text{A11b})$$

so as to absorb the dilution terms. Indeed, Eqs. (A1a)–(A1d) can be reexpressed in terms of the variables above as follows:

$$Hf' = -(w_{\text{rh}} + 1)\hat{\Gamma}_{\delta\phi}f, \quad (\text{A12a})$$

$$Hf'_R = (w_{\text{rh}} + 1)\hat{\Gamma}_{\delta\phi}fR^{1-3w_{\text{rh}}}, \quad (\text{A12b})$$



Sec.	τ_{rh}	T_{rh} (EeV)	$H(T_{\text{rh}})/\hat{\Gamma}_{\delta\phi}$	Y_B (10^{-11})	$Y_{3/2}$ (10^{-13})
A 1	14.3	0.5	0.89	8.55	1.2
A 2	14.5	0.57	0.87	9.6	0.51

FIG. 4. The evolution of the quantities $\log \rho$ (gray dashed line), $\log \rho_R$ (gray line), $\log |Y_L|$ (black solid line), and $\log Y_{3/2}$ (black dashed line) as functions of $\log T$ for the inputs in Eq. (A14). The values of the resulting parameters found by employing our numerical or analytic formulas presented in Appendixes A 1 and A 2, respectively, are listed in the table.

$$Hf'_L = 2\varepsilon_L(w_{\text{rh}} + 1)\hat{\Gamma}_{\delta\phi \rightarrow N_i^c} f R^{3w_{\text{rh}}} / \hat{m}_{\delta\phi}, \quad (\text{A12c})$$

$$Hf'_{3/2} = C_{3/2} n_{\text{eq}}^2 R^3, \quad (\text{A12d})$$

where prime denotes derivation with respect to the logarithmic time

$$\tau = \ln(R/R_i) \Rightarrow R' = R \quad \text{and} \quad R = R_i e^\tau, \quad (\text{A13})$$

with R_i corresponding to the initialization of the $\widehat{\delta\phi}$ oscillations. It can be conveniently selected so that the resolution of the system from $\tau_i = 0$ to $\tau_f > \tau_{\text{rh}}$ is numerically stable—here τ_{rh} corresponds to T_{rh} .

In Fig. 4, we depict the evolution of the quantities $\log \rho$ (dashed gray line), $\log \rho_R$ (solid gray line), $\log |Y_L|$ (black solid line), and $\log Y_{3/2}$ (black dashed line) as functions of $\log T/\text{PeV}$ for a representative set of values of the parameters. In particular, we take

$$N = 0.3, \quad p = 2, \quad \mu/m_{3/2} = 0.1, \quad M_{1N^c} = 33.5 \text{ EeV}. \quad (\text{A14})$$

We use also $N_{\text{st}} = 1/3$ as in Fig. 3. We obtain as outputs

$$N_\star = 55.9, \quad \lambda = 4.8 \times 10^{-5}, \quad \hat{m}_{\delta\phi} = 1.62 \text{ ZeV}, \quad w_{\text{rh}} = 0.28, \quad (\text{A15a})$$

with the rightmost condition in Eq. (19) being fulfilled for $\phi_\star = 0.9315 m_{\text{P}}$. The resulting inflationary observables are

$$n_s = 0.9743, \quad a_s = -4.7 \times 10^{-4} \quad \text{and} \quad r = 0.014, \quad (\text{A15b})$$

whereas the outputs of the postinflationary stage of our scenario are listed in the table of Fig. 4. We there arrange values for various quantities derived by our numerical approach described in this section and by our analytical approximate expressions presented in the next one. From the upper plot of Fig. 4, we observe that $p\text{HI}$ is followed by a ρ -dominated era, due to the oscillating and decaying $\widehat{\delta\phi}$, which lasts at the intersection of the solid and dashed gray lines where T_{rh} is defined. After it, the Universe enters a conventional radiation-dominated era. For $T = T_{\text{rh}}$, the black and dashed solid lines approach their plateau (present) values as shown in the lower panel of Fig. 4.

2. Approximate results

We focus on the regime with $T \geq T_{\text{rh}}$ where $\rho \geq \rho_R$ and, therefore, H in Eq. (A6) assumes the form

$$H \simeq \sqrt{\rho}/\sqrt{3}m_{\text{P}}. \quad (\text{A16})$$

Inserting it into Eq. (A12) we find

$$\rho \simeq \rho_i e^{-3(w_{\text{rh}}+1)\tau} \quad \text{with} \quad \rho_i = \rho(0) \quad (\text{A17})$$

given in Eq. (A8). Upon substitution in Eq. (A12b), we obtain

$$\rho_R(\tau) = \rho_{\text{Ri}}(e^{-3(1+w_{\text{rh}})\tau/2} - e^{-4\tau}), \quad (\text{A18})$$

where we define ρ_{Ri} as follows:

$$\rho_{\text{Ri}} = \frac{2\sqrt{3}(w_{\text{rh}}+1)}{5-3w_{\text{rh}}} m_{\text{P}} \hat{\Gamma}_{\delta\phi} \rho_i^{1/2}. \quad (\text{A19})$$

The function $\rho_R(\tau)$ in Eq. (A18) reaches a maximum at

$$\tau_{\text{max}} \simeq \frac{2}{5-3w_{\text{rh}}} \ln \frac{8}{3(w_{\text{rh}}+1)}, \quad (\text{A20})$$

which turns out to be very close to $\tau = \tau_i$, e.g., $\tau_{\text{max}} - \tau_i = 0.39$ or 0.35 for $w_{\text{rh}} = 0$ or $w_{\text{rh}} = 1/3$, respectively. Therefore, for $\tau > \tau_{\text{max}}$ and $w_{\text{rh}} < 5/3$, ρ_R in Eq. (A18) is dominated by the first term in the parentheses. Upon substitution into the rightmost relation in Eq. (A7), we find that T decreases as follows:

$$T = T_i e^{-3(w_{\text{rh}}+1)\tau/8} \quad \text{with} \quad T_i = \left(\frac{30\rho_{\text{Ri}}}{g_\star \pi^2} \right)^{1/4}. \quad (\text{A21})$$

For simplicity, we take a constant value for $g_\star$ throughout.

a. Reheat temperature

The completion of the reheating takes place when Eq. (A9) is satisfied. Taking into account Eqs. (A17) and (A18), we find that τ_{rh} takes the form

$$\tau_{\text{rh}} \simeq \frac{2}{3(w_{\text{rh}} + 1)} \ln \frac{\rho_i}{\rho_{\text{Ri}}}. \quad (\text{A22})$$

If we insert this result into Eq. (A21) and make use of Eq. (A19), we arrive at the result of Eq. (40) which does not depend on ρ_i . Our approach is consistent with an alternative definition of T_{rh} according to which it is defined from the condition $H(T_{\text{rh}}) \simeq \hat{\Gamma}_{\delta\phi}$. Indeed, we find

$$H(\tau_{\text{rh}})/\hat{\Gamma}_{\delta\phi} = 2\sqrt{2}(w_{\text{rh}} + 1)/(5 - 3w_{\text{rh}}), \quad (\text{A23})$$

with numerical values given in the table of Fig. 4.

b. Lepton-asymmetry abundance

Integrating Eq. (A12c) with H given by Eq. (A16) we find

$$n_L = \frac{4(w_{\text{rh}} + 1)\epsilon_L m_{\text{P}}}{\sqrt{3}(1 - w_{\text{rh}})\hat{m}_{\delta\phi}} \rho_i^{1/2} \hat{\Gamma}_{\delta\phi} (e^{-3(w_{\text{rh}}+1)\tau/2} - e^{-3\tau}), \quad (\text{A24})$$

where the first exponential in the parentheses dominates over the latter for $w_{\text{rh}} < 1$. Confining ourselves to this regime and applying the expression above for $\tau = \tau_{\text{rh}}$ given by Eq. (A22), we obtain

$$n_L(\tau_{\text{rh}}) = \frac{8(w_{\text{rh}} + 1)^2 \epsilon_L m_{\text{P}}^2}{(1 - w_{\text{rh}})(5 - 3w_{\text{rh}})\hat{m}_{\delta\phi}} \frac{\hat{\Gamma}_{\delta\phi \rightarrow N_i^c}}{\hat{\Gamma}_{\delta\phi}} \hat{\Gamma}_{\delta\phi}^2, \quad (\text{A25})$$

where ρ_i is again eliminated if we take into account Eq. (A19). Plugging into the expression above $\hat{\Gamma}_{\delta\phi}^2$ from Eq. (40) and employing the rightmost expression in Eq. (A7), we end up with the result of Eq. (41) after

including the numerical factor (-0.35) which accounts for the sphaleron effects.

c. Gravitino abundance

Solving Eq. (A12d) with H and T given by Eqs. (A16) and (A21), respectively, we find

$$n_{3/2} = \frac{4\zeta(3)^2 m_{\text{P}} C_{3/2} T_i^6}{\sqrt{3}\pi^4 \rho_i^{1/2} (3 - w_{\text{rh}})} (e^{-3(w_{\text{rh}}+1)\tau/4} - e^{-3\tau}). \quad (\text{A26})$$

The first exponential in the parentheses above is dominant for $w_{\text{rh}} < 3$. Working in that range of w_{rh} values and replacing $\tau = \tau_{\text{rh}}$ from Eq. (A22), we arrive at

$$n_{3/2}(\tau_{\text{rh}}) = \left(\frac{10}{g_*}\right)^{1/2} \frac{4\zeta(3)^2 m_{\text{P}} C_{3/2}}{\pi^5 (3 - w_{\text{rh}})} T_{\text{rh}}^4, \quad (\text{A27})$$

where ρ_i is again eliminated, as can be shown by replacing T_{rh}^4 and ρ_{Ri} from Eqs. (40) and (A19), respectively. Employing the rightmost expression in Eq. (A7), we end up with

$$Y_{3/2} = \frac{n_{3/2}}{s}(\tau_{\text{rh}}) = \frac{90}{g_*} \left(\frac{10}{g_*}\right)^{1/2} \frac{\zeta(3)^2 C_{3/2}}{\pi^7 (3 - w_{\text{rh}})} m_{\text{P}} T_{\text{rh}}. \quad (\text{A28})$$

Comparing the result above with the numerical one—see the table of Fig. 4—we remark that there is a slight discrepancy, which turns out to be systematic. Its origin stems from the fact that there is a residual production of $\tilde{G}$'s for $\tau > \tau_{\text{rh}}$. To account for this effect, we multiply the result of Eq. (A28) by a factor of 2 and include it in the numerical factor of Eq. (45), achieving, thereby, a better agreement with numerics.

In conclusion, we motivated the adopted formulas in Eqs. (40), (41), and (45), which exhibit an explicit dependence on w_{rh} and allow us to analyze consistently the postinflationary dynamics. The numerical impact of this improvement on the results is a correction of order unity and, therefore, it is not so crucial for the viability of our setup.

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