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# The inflaton as an MSSM Higgs and open string modulus monodromy inflation

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ABSTRACT

It has been recently pointed out that high scale inflation, as recently hinted by the BICEP2 results, is consistent with the identification of an inflaton mass $m_I \simeq 10^{13}$ GeV with the SUSY breaking scale in an MSSM with a fine-tuned SM Higgs. This identification leads to a Higgs mass $m_h \simeq 126$ GeV, consistent with LHC measurements. Here we propose that this naturally suggests to identify the inflaton with the heavy MSSM Higgs system. The fact that the extrapolated Higgs coupling $\lambda_{SM} \simeq 0$ at scales below the Planck scale suggests the Higgs degrees of freedom could be associated with a Wilson line or D-brane position modulus in string theory. The Higgs system then has a shift symmetry and an $N = 2$ structure which guarantees that its potential has an approximate quadratic chaotic inflation form. These moduli in string compactifications, being compact, allow for trans-Planckian inflaton field range analogous to a version of monodromy inflation.



## 1. Introduction

Inflationary cosmology has achieved an impressive series of successful tests. The recent polarisation measurements by the BICEP2 Collaboration [1] are consistent with the presence of cosmological tensor fluctuations in the very early universe. If the cosmological origin of these tensor polarisations is confirmed, it is intriguing that large field inflationary models and in particular the chaotic inflationary scenario with a simple quadratic inflation potential [2] would possibly provide for the simplest explanation. On the other hand the CMS and ATLAS Collaborations confirmed the existence of a Higgs boson with mass around 126 GeV [3]. It has always been tempting to identify the inflaton with the Higgs boson since, after all, the Higgs boson is the only known fundamental scalar which has been observed. It was however soon found that the form of the SM potential, with a quartic term and a small mass parameter was not appropriate to generate successful inflation. Modifications were proposed using non-minimal couplings of the Higgs boson to the curvature, leading to inflationary models with an effective Starobinsky like structure [4]. These models however are not free of problems (see e.g. [5] and references therein) and moreover the LHC measurements of the Higgs mass do not favour them. This is because, for the observed value of the Higgs and top quark masses, an RGE extrapolation of the Higgs quartic coupling $\lambda_{SM}$ shows that the SM potential becomes unstable at energies of order $10^{11}$–$10^{13}$ GeV, well below the Planck scale [6]. Furthermore such models predict very small tensor fluctuations, in apparent contradiction to BICEP2 results.

The fact that the SM potential becomes unstable at a scale $10^{10}$–$10^{13}$ GeV suggests that at those scales some new physics sets on stabilising the potential. In [7,8] it was suggested that that scale could correspond to the SUSY breaking scale and in [9] it was found that indeed that assumption is consistent with the observed Higgs mass (see also [10,11]). If this is the case, the role of SUSY would not be stabilising the Higgs mass, which would have to be fine-tuned [12], but rather to stabilise the potential. From the point of view of string theory, the existence of SUSY at some scale, not necessarily the TeV scale, is strongly motivated, since it is a built-in symmetry of the theory and provides stability for the abundant scalars appearing in string compactifications. Furthermore, the fine-tuning of the Higgs mass could be motivated from the point of view of the string landscape.

In Ref. [13] we proposed that the polarisation BICEP2 results, if indeed pointing to cosmological tensor perturbations, are consistent with the identification of an inflaton mass $m_I \simeq 10^{13}$ GeV with the SUSY breaking scale in an MSSM with a fine-tuned SM Higgs. We showed how this identification led to results for the Higgs mass consistent with the LHC results. Here we show that, if indeed an MSSM-like structure is realised at an intermediate scale $M_{ss} \simeq 10^{13}$ GeV, the MSSM Higgses $h$, $H$ can give rise to inflation. We show that a quadratic mass term leading to a chaotic inflation naturally appears if an appropriate symmetry structure is present in the Higgs sector. We also show that these kinds of

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symmetry appear in string compactifications and higher dimensional models. Large trans-Planckian inflaton range may appear in a way analogous to that of monodromy inflation.

## 2. The intermediate scale MSSM and inflation

The minimal Higgs system in the MSSM has two EW doublets $H_u$ and $H_d$. The scalar potential for the neutral scalars is given by a D-term and general SUSY-breaking soft terms, with a general structure

$$V_{Higgs} = m_u^2|H_u|^2 + m_d^2|H_d|^2 + (BH_uH_d + \text{h.c.}) + \frac{g^2+g_1^3}{8}\left(|H_u|^2 - |H_d|^2\right)^2 \tag{1}$$

where $m_{u,d}^2$ includes both the soft masses and a possible contribution of a SUSY $\mu$-term. All of them will be of order $10^{13}$ GeV, and a massless SM Higgs doublet would result from a delicate fine tuning of the mass parameters [7–9]. Here we will concentrate in the two complex neutral scalars. Let us define the two eigenvalues of the mass matrix as

$$h = \sin\beta H_u - \cos\beta H_d^*; \qquad H = \cos\beta H_u + \sin\beta H_d^* \tag{2}$$

with respective masses $m_-^2$ and $m_+^2$ given by

$$m_\pm^2 = \frac{1}{2}\left(m_u^2 + m_d^2 \pm \sqrt{\left(m_u^2 - m_d^2\right)^2 + 4|B|^2}\right) \tag{3}$$

Note that a zero eigenvalue, corresponding to $m_h = m_- = 0$ appears when $|B|^2 = m_u^2 m_d^2$, yielding an (approximately) massless Higgs $h$. This we want to happen *at the SUSY-breaking scale $M_{ss}$*. Note however that, running-up in energies to the GUT scale $m_-^2$ will be positive, and both $m_\pm^2$ will be not vanishing at the GUT scale. Without loss of generality we can take the neutral vevs of $h$, $H$ real. We then get

$$V_{Higgs} = m_-^2 h^2 + m_+^2 H^2 + \frac{g^2+g_1^2}{8}\left(\cos 2\beta\left(H^2 - h^2\right) + 2hH\sin 2\beta\right)^2. \tag{4}$$

Note that, close below the $M_{ss}$ scale where $m_+ \gg m_-$, one recovers a SM Higgs potential with

$$V_{SM} = m_-^2 h^2 + \frac{g^2+g_1^2}{8}\cos^2 2\beta|h|^4. \tag{5}$$

As we said, at such high scales we know that the Higgs self-coupling $\lambda \simeq 0$, which in the present context implies that $\cos^2 2\beta \simeq 0$ at $M_{ss}$, recovering the results in [7–10]. At scales $\gtrsim M_{ss}$ the scalar potential is then given by

$$V_{Higgs} = m_-^2 h^2 + m_+^2 H^2 + \frac{g^2+g_1^2}{2}\left(h^2H^2\right), \tag{6}$$

with $m_-^2 \lesssim m_+^2$ and $\cos^2 2\beta \simeq 0$, as suggested by the low-energy Higgs mass results. Note that along the direction $h = 0$, corresponding to a very small vev for the SM Higgs, the other MSSM Higgs scalar $H$ has a chaotic inflation scalar potential. The inflaton/Higgs potential starts with very large vevs as in conventional chaotic inflation. As the inflation goes down eventually $H$ finds a minimum at $H = 0$, forced by the large mass $m_+$ term present at $M_{ss}$, and oscillate, reheating the universe. Note that the reheating proceeds dominantly through SM particles. The SM Higgs $h$ has a fine-tuned mass $m_-^2$ which is (approximately) zero around $M_{ss}$ (although is positive and of order $M_{ss}$ at larger scales, due to RGE running).

## 3. Mass scales and string theory

For this system to work we have to check for the stability of the Higgs/inflaton potential. Furthermore we know that slow-roll and large tensor perturbations suggest the inflaton field should ride along trans-Planckian regions. Finally, we would like to know what is the origin of the mass scale $m \simeq 10^{13}$ GeV which fixes both the SUSY breaking scale and the inflaton mass. To answer all these questions we need an UV completion of the theory which in what follows we assume to be string theory.

First let us discuss the origin of the SUSY breaking scale. The natural option in string theory is to consider the effect of antisymmetric closed string fluxes, which for an arbitrary choice lead to SUSY breaking masses. These are particularly well understood in the case of Type IIB orientifold compactifications. Consider for example a D7-brane wrapping a 4-cycle $\Sigma$ in a compact Calabi–Yau (CY) space. The DBI action for the brane has the general form [14]

$$S = -\frac{1}{g_s(\alpha')^4(2\pi)^7}\int_\Sigma d^8x\sqrt{-\det\left(P[G+B] - 2\pi\alpha' F\right)}, \tag{7}$$

where $g_s$ is the string coupling, $\alpha'$ is the (inverse) string tension, and $F$ is the Yang–Mills field strength. $P[G]$ is the induced metric and $P[B]$ is the pull-back of the antisymmetric $B_{ij}$ NS field. One can integrate locally for the $B$-field in terms of its field strength 3-form $H_3$, $B_{ij} = H_{ijk}z^k$, with $z_k$ a coordinate in the CY transverse to $\Sigma$. This coordinate is parametrised by the vev of a scalar $\Phi$, so that one has $B \simeq H_3\Phi$. For non-vanishing fluxes $\langle H_3\rangle \neq 0$, the expansion of the action for diluted fluxes induces a mass term of the form $\langle H_3^2\rangle|\Phi|^2$. The fluxes are Dirac-quantised as $\int_\gamma H_3 = (2\pi)^2\alpha' n_\gamma$, with $\gamma$ a 3-cycle in the CY and $n_\gamma$ an integer. Thus for an isotropic compactification one expects $H_3 \simeq \alpha'/R^3$, with $R^6$ the CY volume and hence one gets

$$m_\Phi^2 \simeq H_3^2 \simeq \frac{(\alpha')2}{R^6} \simeq \frac{M_s^4}{M_p^2}, \tag{8}$$

where $M_s^2 = (\alpha')^{-1}$, the string scale. Taking the string scale to be of order the unification scale $M_s \simeq 10^{16}$ GeV, one obtains soft terms of the required size, $m_\Phi \simeq 10^{13}$ GeV. So the generic presence of antisymmetric fluxes in string theory would provide for an explanation for a SUSY breaking scale of that size. Note that for a given compactification there is a variety of fluxes (NS, RR, non-geometric, ...) that may be turned on, leading to a variety of soft terms, see e.g. [14–16]. Some particular classes of fluxes may also give rise to supersymmetric couplings, like a $\mu$-term. All in all we have a hierarchy of mass scales

$$M_{ss} \simeq 10^{13}\ \text{GeV} < M_c, \qquad M_s \simeq 10^{16}\ \text{GeV} < M_p. \tag{9}$$

Here $M_c$ is the compactification scale that, e.g. in this Type IIB setting is given by $M_c \simeq M_s(\alpha_{GUT}/2g_s)^{1/4}$, and hence is only slightly below the string scale. Note in particular that using a field theory scalar potential above the unification scale $10^{16}$ GeV is questionable. We discuss this point below. In addition, if the closed string moduli are also fixed by fluxes, one should include them in the full scalar potential. In what follows we assume that the moduli are fixed at a higher scale than $M_{ss}$ respecting SUSY, so that one can consistently focus on the inflaton/Higgs potential. One could obtain such a separation of scales with appropriate flux and geometry choices. In any event that would be very model dependent and we leave for future investigation.

## 4. Trans-Planckian inflaton

The large tensor perturbations detected at BICEP2 suggest a trans-Planckian field range for the inflaton [17]. On the other hand, as we said, using a field theory potential is questionable for field vevs above the compactification/string scales $M_c \simeq M_s \simeq 10^{16}$ GeV. A very elegant solution to this general problem was suggested in [18] (see also [19,20]). If the inflation is identified with an axion-like field $a$ with a classical shift invariance $a \to a + c$, with a non-trivial monodromy field space, large inflation values may be achieved without trans-Planckian axion decay constants. In other words, the inflation range is not directly limited by the size of the manifold. Interestingly enough, this idea also applies to Wilson line and D-brane position moduli in string theory, which also present shift symmetries, in a variety of cases.

This suggests to consider the Higgs sector as associated with Wilson lines in string compactifications, to make the large field limit consistent in this set-up. In fact, as emphasised in [10], it is an intriguing fact that the apparent vanishing of the Higgs self-coupling at scales above $10^{11}$ GeV may be understood in terms of a shift invariance

$$H_{u,d} \to H_{u,d} + i\sigma \tag{10}$$

in the quadratic potential (here $\sigma$ is real). Indeed, the quadratic potential is only invariant if $\cos^2 2\beta = 0$. Under this shift the $h$ and $H$ fields transform as

$$\begin{aligned} h &\to h + i\sigma(\sin\beta + \cos\beta), \\ H &\to H - i\sigma(\sin\beta - \cos\beta). \end{aligned} \tag{11}$$

Then for $\tan\beta = 1$, $h \to h + i\sqrt{2}\sigma$ and $H \to H$, and a shift symmetry appears for the $h$ field. The field $h$ in this limit is massless, $m_- = 0$ but $m_+ \neq 0$. This symmetry would not be exact in the MSSM, since e.g. loop corrections involving Yukawa couplings affect differently the $H_u$ and the $H_d$ masses, but still, the fact that $\lambda_{SM} \simeq 0$ at large scales may be taken as an indication of an approximate shift invariance at some large scale.

Analogous shift symmetries are known to be present in certain subsectors of string compactifications [22]. In particular, in heterotic $Z_{2N}$ toroidal orbifold compactifications, the untwisted charged fields $H_1$, $H_2$ associated with complex planes with a twist of order two have a Kahler potential

$$\kappa_4 K = -\log\big((U+U^*)(T+T^*) - \alpha'(H_1 + H_2^*)(H_1^* + H_2)\big),$$

where $U$ and $T$ are the complex structure and Kahler modulus of the $T^2$ torus associated with the mentioned complex plane. Note that the Kahler potential is invariant under a shift symmetry $H_{1,2} \to H_{1,2} + i\sigma$. As noted in [22], if SUSY-breaking is induced by the auxiliary fields of the moduli or the dilaton (no matter what combination), the quadratic part of the scalar potential may be written as

$$V \propto (H_1 + H_2^*)(H_1^* + H_2), \tag{12}$$

which is explicitly invariant under the shift symmetry, and would correspond to a mass term $m_+^2|H|^2$ and $m_-^2 = 0$ in the Higgs case. In the heterotic case this shift symmetry is a remnant of the gauge transformations of a 6D gauge boson, and the matter fields correspond to a continuous Wilson-line moduli. This shift symmetry has also been exploited in the context of models with extra dimensions under the name of 'gauge-Higgs unification', see [23]. An additional important ingredient in these string theory settings is that the $H_{1,2}$ fields appear in a subsector of the theory respecting $N = 2$ supersymmetry, with them forming an $N = 2$ hypermultiplet. E.g. in the heterotic orbifold examples this happens because the fields $H_{1,2}$ appear from an $N = 2$ sector of the compactification associated with a complex plane with an order-2 twist. This extended supersymmetry forbids then the appearance of any dim $= 4$ operator F-term contribution to the scalar potential involving just the $H_{1,2}$ fields.

Moduli fixing and SUSY-breaking induced by fluxes is better understood in the context of Type II orientifolds. Shift symmetries for D-brane positions and/or Wilson line open string moduli also appear in Type II string constructions, as expected from string dualities. Indeed, by S-duality on recovers the same structure of Wilson line moduli in Type IIB orientifolds with D9-branes. Further T-dualities yield orientifolds with matter fields living on D3-branes and or D7-branes, and Wilson lines mapping to either Wilson line moduli or D-brane position moduli. The latter could perhaps be the simplest way in trying to implement the idea of monodromy for a Higgs inflation. Recently it has been shown in [20] how Wilson-line monodromy inflation may be quite generic in Type II orientifold models with $Dp$ branes wrapping $(p-3)$-cycles in a CY (see also [21,24]). In the simplest implementation one can summarise the idea by saying that any periodic string moduli, either an axion, D-brane position moduli or Wilson line moduli give rise to a monodromy potential in the presence of different types of closed string fluxes. For a recent discussion with the D7-brane position acting as an inflaton see [21]. Let us give a simple MSSM-like toy model using an example with D7-branes on $Z_N$ singularities.

Consider a set of 6 D7-branes wrapping a $T^4$ in a CY with local geometry $(\mathbf{T^4} \times \mathbf{C})/Z_4$ and located on a $\mathbf{Z_4}$ singularity, with local coordinates twisted by $(z_1, z_2, z_3) \to (\alpha z_1, \alpha z_2, \alpha^2 z_3)$, with $\alpha = \exp(i2\pi/4)$. On the open strings there is a Chan–Paton matrix $\gamma = \mathrm{diag}(\alpha\mathbf{1_3}, \alpha^2\mathbf{1_2}, \mathbf{1})$, with $\alpha = \exp(i2\pi/4)$. The open string sector includes gauge bosons in the gauge group $U(3) \times U(2) \times U(1)$ with matter fields transforming like (see e.g. [15])

$$2(3, \bar{2}) + 2(\bar{3}, 1) + (1, 2) + (1, \bar{2}).$$

To get RR-tadpole cancellation there must be additional D3-branes at the singularities. The open D3–D7 open strings complete the spectrum to two generations of the SM plus extra vector-like matter which is not relevant for the discussion. The main point is that, associated with the 3-d complex plane which suffers a $Z_2$ twist, there is a vector-like set of Higgs multiplets $(1,2) + (1,\bar{2})$ which may be identified with $H_u$, $H_d$. The vev of $(H_u + H_d^*)$ parametrises the location of the D7-branes in the $z_3$ coordinate. We will assume that this vev has a periodic behaviour around a 1-cycle in compact dimensions. E.g., one may consider instead a local structure $\mathbf{C^2} \times \mathbf{T^2}$ with $z_3$ living in the 2-torus. One can then consider the addition of antisymmetric RR and NS IIB closed string fluxes, as we discussed above (one has to consider also contributions from the D7-brane Chern–Simons action). In the presence of ISD $G_3$ fluxes with non-SUSY $G_{(0,3)}$ and SUSY $S_{\bar{3}\bar{3}}$ components, soft terms are induced, as we mentioned above. In this example one gets soft terms of the form [15]

$$M_3 = M_2 = M = \frac{g_s^{1/2}}{3\sqrt{2}} G^*_{(0,3)}$$

$$m^2_{H_u} = m^2_{H_d} = |M|^2, \qquad \mu = -\frac{g_s^{1/2}}{6\sqrt{2}} S^*_{\bar{3}\bar{3}}, \qquad B = M\mu$$

where $G$ and $S$ are flux densities at the singularity, $M$ is the gaugino masses, $\mu$ is a SUSY mass term for the Higgs and $B$ is the standard Higgs scalar bilinear term. These mass terms provide for a specific origin for the MSSM scalar potential of Eq. (1), with the additional ingredient that the Higgs fields may get large field values through the monodromy induced by the fluxes, which break the shift symmetry. Note that fluxes may also induce a SUSY $\mu$-term, and that at the string/unification scale $\det(m^2_{Higgs}) > 0$, and there is no zero eigenvalue. Such zero eigenvalue, giving rise to a light SM Higgs, may arise however at lower energies $\simeq 10^{13}$ GeV upon RGE running [7–9]. In the scalar potential there is a exchange symmetry $H_u \leftrightarrow H_d$ which will be eventually broken by the stronger Yukawa coupling of $H_u$ to the top quark mass, driving $m_u^2 < m_d^2$ [8,9]. Although this example is not fully realistic, it illustrates how the required monodromy for the Higgs fields may easily appear in Type IIB models with closed string fluxes. As we already said, a final point to remark is that in these schemes the issue of inflation potential and the fixing of the closed string moduli are necessarily interrelated. Still one can play with the different volumes and 1-cycle sizes so that the moduli are fixed at scales larger than the inflaton mass [20]. Thus a fully consistent model should also include an appropriate treatment of the moduli fixing. Furthermore, for large inflation values the behaviour of the scalar potential may be modified, depending on the particular geometric implementation of the monodromy, see [18,20].

## 5. Discussion

In this note we have proposed that the SM Higgs boson and the inflaton are SUSY partners within an MSSM structure at a high SUSY breaking scale. This leads to identify the inflaton with the heavy scalar Higgs field $H$ which is present in the MSSM in addition to the standard model Higgs $h$. For this to be the case the SUSY-breaking scale should coincide with the inflaton mass $m \simeq 10^{13}$ GeV, as recently suggested in [13]. Within the context of string theory, such a scale naturally appears since for a string scale $\simeq 10^{16}$ GeV, flux-induced soft terms are of order $M_{ss} \simeq M_s^2/M_p \simeq 10^{13}$ GeV. The induced mass terms give rise to a chaotic inflationary model. The fact that the Higgs self-coupling $\lambda_{SM}$ seems to vanish at an intermediate scale $\simeq 10^{11}$–$10^{13}$ GeV, suggests the existence of an approximate shift symmetry in the MSSM Higgs system. Such type of symmetries are characteristic of open string moduli in Type II string compactifications. It has been recently realised [20] that open string moduli, in the presence of appropriate closed string fluxes, naturally give rise to a simple version of monodromy inflation (see also [21]). Large inflaton field values, as required by large field inflation and large tensor fluctuations, naturally appear in these schemes. It would be very interesting to obtain more complete string compactifications in which the Higgs doublets may be associated with a D-brane position/Wilson line moduli with non-trivial monodromy as in the toy models here suggested. Work along these lines is in progress.

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