Evidence for inflation in an Axion Landscape

PASCOS2018, Cleveland, June 4-8, 2018

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June 8, 2018

Outline

- Axion inflation
- A stronger case for SUSY post Higgs boson discovery

Inflation

- Inflationary models resolve a number of problems associated with Big Bang cosmology which include the flatness problem, the horizon problem, and the monopole problem [1].
- The astrophysical data from the Planck experiment [2] has put significant constraints on models eliminating some and reducing the parameter space of others.
- Although there is a large amount of work on inflation, there is currently no standard model of inflation.
- I will discuss a recently proposed axion model [3] within supersymmetry and supergravity which resolves problems of axion decay constant of previous axion models and is consistent with Planck data.

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1 Guth:1980,Starobinsky:1980,Linde:1981,Sato81,Albrecht and Steinhardt:1982, Linde:1983

2 P. A. R. Ade *et al.* [Planck Collaboration], Astron. Astrophys. **594**, A20 (2016) [arXiv:1502.02114 [astro-ph.CO]].

3 P,N., M. Piskunov, JHEP **1803**, 121 (2018) doi:10.1007/JHEP03(2018)121 [arXiv:1712.01357 [hep-ph]].

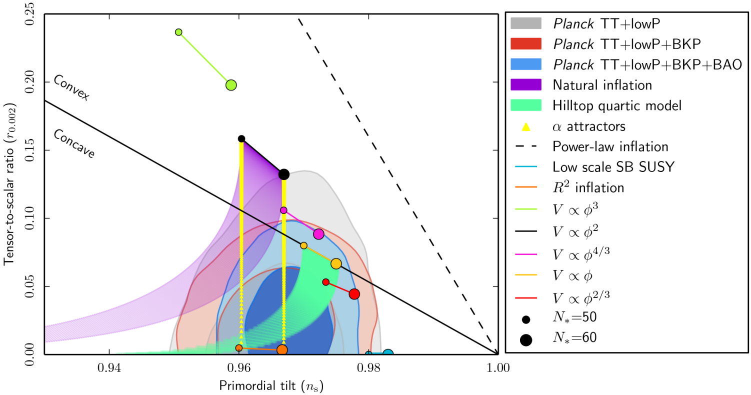


Figure from P. A. R. Ade *et al.* [Planck Collaboration], Astron. Astrophys. **594**, A20 (2016) [arXiv:1502.02114 [astro-ph.CO]].

Axionic inflation

An early work using the QCD axion is the so -called natural inflation where the inflation potential is of the form[4]

$$V(a) = \Lambda^4 \left(1 + \cos(\frac{a}{f})\right) ,$$

f is axion decay constant.

- For a QCD axion $10^9 < f < 10^{12}$ GeV. For inflation one requires $f > 10M_{Pl}$.
- $f > M_{Pl}$ is undesirable since global symmetry is not preserved by quantum gravity unless it has a gauge origin.
- Further, string theory prefers f in the range $(10^{16} - 10^{18})$ GeV.

[4] K. Freese, J. A. Frieman and A. V. Olinto, Phys. Rev. Lett. **65**, 3233 (1990).

Alignment mechanism [5]

One suggestion to realize $f < M_{Pl}$ is to use two axions and the alignment mechanism to achieve a flat direction. For a model with two axions ϕ_1 and ϕ_2 one considers a potential

$$V(\phi) = \Lambda_1^4 \left[1 - \cos\left(\frac{\phi_1}{f_1} + \frac{\phi_2}{f_2}\right)\right] + \Lambda_2^4 \left[1 - \cos\left(\frac{\phi_1}{f_3} + \frac{\phi_2}{f_4}\right)\right] . \tag{1}$$

- Constrain for a flat direction:
$$\frac{f_1}{f_2} = \frac{f_3}{f_4}.$$
- For a recent attempt to implement it in string theory see[6].

[5] J. E. Kim, H. P. Nilles and M. Peloso, JCAP **0501**, 005 (2005).

[6] C. Long, L. McAllister and P. McGuirk, Phys. Rev. D **90**, 023501 (2014).

- We propose an alternative to alignment mechanism which is a Fast-Roll-Slow -Roll classification of fields in an axion landscape. Inflation is governed by the slow roll field. This technique is general and applicable to multi-field inflation with a $U(1)$ symmetry.

  We will discuss explicit models to implement Fast-Roll-Slow -Roll decomposition of the axion potential.

  Specifically

- Inflation in an axion landscape in SUSY/Supergravity

- Axion inflation in SUSY Dirac-Born-Infeld type framework and non-Gaussianity

$U(1)$ symmetry and Inflation in an Axion Landscape[7]

- We consider a landscape of chiral fields charged under a global $U(1)$ symmetry. We will call the real parts of chiral fields saxions and the imaginary parts axions. There is only one pseudo - Nambu - Goldstone - Boson (pNBG) which will act as the inflaton.

- We consider the case when the pNGB is not the axion of QCD but rather one that comes from strings.

- The effective low energy theory for the pNGB is very different from the old natural inflation models and one can generate inflation with $f < M_{Pl}$.

[7] P,N., Maksim Piskunov, JHEP **1803**, 121 (2018) doi:10.1007/JHEP03(2018)121 [arXiv:1712.01357 [hep-ph]].

String Axions

- An nice discussion of axions in strings is given in

P. Svrcek and E. Witten, "Axions In String Theory," JHEP **0606**, 051 (2006).

- String axions have recently been used in models of ultra-light or fuzzy dark matter

  - L. Hui, J. P. Ostriker, S. Tremaine and E. Witten, "Ultralight scalars as cosmological dark matter," Phys. Rev. D **95**, no. 4, 043541 (2017).
  - J. E. Kim and D. J. E. Marsh "An ultralight pseudoscalar boson," Phys. Rev. D **93**, no. 2, 025027 (2016)
    vs
  - J. Halverson, C. Long and P.N., "Ultralight axion in supersymmetry and strings and cosmology at small scales," Phys. Rev. D **96**, no. 5, 056025 (2017)

The Model

- Suppose we have a set of fields Φ_i $(i = 1, \cdots, m)$ where Φ_i carry the same charge under the shift symmetry and the fields $\bar{\Phi}_i$ $(i = 1, \cdots, m)$ carry the opposite charge.
- We may parametrize ϕ_k and $\bar{\phi}_k$ so that

$$\phi_k = (f_k + \rho_k) e^{i a_k / f_k}, \quad \bar{\phi}_k = (\bar{f}_k + \bar{\rho}_k) e^{i \bar{a}_k / \bar{f}_k},$$

- This allows us to write a non-trivial superpotential which can stabilize the saxions.

$$W = W_s + W_{sb}$$
$$W_{sb} = \sum_{r=1}^{q} (P_r(\Phi) + \bar{P}_r(\bar{\Phi}))$$

W_s is symmetry preserving and W_{sb} breaks the shift symmetry. $P_r(\Phi), \bar{P}_r(\bar{\Phi})$ are polynomials of power r in the fields. Saxions are stabilized via constraints

$$W_{,\phi} = 0 = W_{,\bar{\phi}}.$$

- $2m$ axionic fields: $a_1, \cdots, a_m$ and $\bar{a}_1, \cdots, \bar{a}_m$.
- Fast roll-slow roll basis:
  - Fast roll fields are invariant under the shift symmetry

$$b_k = \frac{a_{k+1}}{f_{k+1}} - \frac{a_1}{f_1}, \quad k = 1, 2, \cdots, m-1\,,$$
$$\bar{b}_k = \frac{\bar{a}_{k+1}}{\bar{f}_{k+1}} - \frac{\bar{a}_1}{\bar{f}_1}, \quad k = 1, 2, \cdots, m-1\,,$$
$$b_+ = \frac{a_1}{f_1} + \frac{\bar{a}_1}{\bar{f}_1}\,.$$

  - b_- is the pNGB, and is the inflaton and the slow roll field

$$b_- = \frac{1}{\sqrt{\sum_{k=1}^m f_k^2 + \sum_{k=1}^m \bar{f}_k^2}} \left(\sum_{k=1}^m f_k a_k - \sum_{k=1}^m \bar{f}_k \bar{a}_k \right)\,.$$

Slow roll potential with stabilized saxions

$$V(b) = V_{\text{fast}} + V_{\text{slow}}(b_-)$$

$$V_{\text{slow}} = \sum_{r=1}^{q} C_r \left(1 - \cos\left(\frac{r}{f_e} b_-\right)\right) + \sum_{s=1}^{q} \sum_{r=s+1}^{q} C_{rs} \left(1 - \cos\left(\frac{r-s}{f_e} b_-\right)\right)$$

where

$$f_e = \sqrt{\sum_{k=1}^{m} f_k^2 + \sum_{k=1}^{m} \bar{f}_k^2}\,.$$

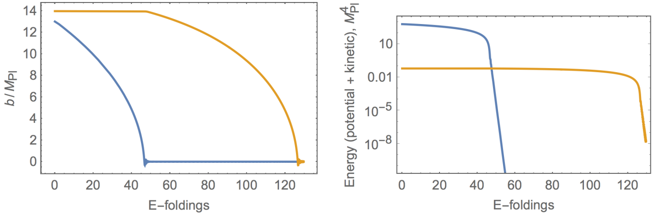
A remarkable aspect of V_{slow} is that it depend only on an effective decay constant f_e. For N number of fields and $f_k = f = \bar{f}_k$,

$$f_e = \sqrt{N} f$$

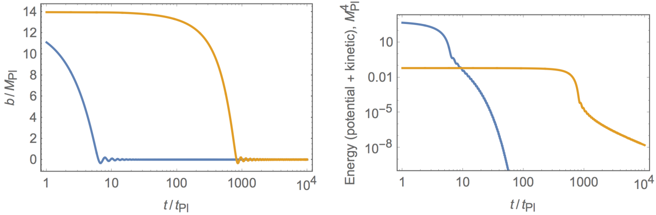
- N-flation emerges in a natural way.
- This is one of the two ways in which $f_e >> M_{Pl}$ can be achieved with sub-Planckian f for an appropriately large N.

The superposition of several cosines give rise to a relatively flat potential for the slow component so that inflation can occur

Fast and slow components vs E-foldings

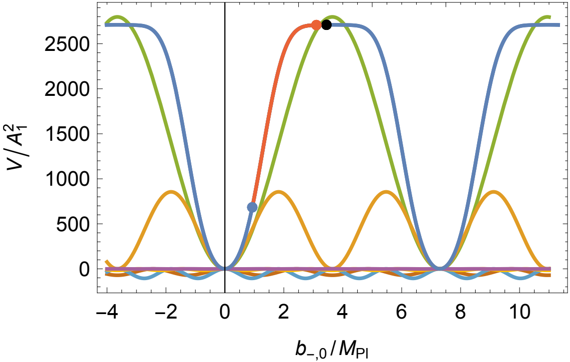


Fast and slow components vs time in Planck units: $\boldsymbol{t_{pl} = 5.4 \times 10^{-44}}$s.



The superposition of several cosines makes the model very different from the natural inflation model. For the case $m = 1$ so we have just one pair of axions and we consider $q = 3$ which gives a superposition of six cosines.

Green curve: Plot of the largest cosine.
Brown curve: Plot of the second largest cosine.
Blue curves+ red curves: Generic inflation potentials.



Emergence of a locally flat potential

- The case $q = 3$.

$$\frac{b_-}{\sqrt{2}f} = \pi : \begin{cases} \text{max} \quad \text{for} \quad 1\text{st}, 3\text{rd}, 5\text{th} \\ \text{min} \quad \text{for} \quad 2\text{nd}, 4\text{th}, 6\text{th} \end{cases} .$$

- V_{slow} for the two axion model with $q = 3$

$$\begin{aligned} V\left(b_-\right) = {} & 1398.96\left(1 - \cos\left(\frac{b_-}{\sqrt{2}f}\right)\right) + 427.466\left(1 - \cos\left(\frac{2b_-}{\sqrt{2}f}\right)\right) \\ & - 35.4939\left(1 - \cos\left(\frac{3b_-}{\sqrt{2}f}\right)\right) - 52.9837\left(1 - \cos\left(\frac{4b_-}{\sqrt{2}f}\right)\right) \\ & - 8.28504\left(1 - \cos\left(\frac{5b_-}{\sqrt{2}f}\right)\right) - 0.632442\left(1 - \cos\left(\frac{6b_-}{\sqrt{2}f}\right)\right) . \end{aligned}$$

$$f_e/f = \frac{\sqrt{\sum_{k-odd} \Lambda_k^4}}{\sqrt{\left[\sum_{k-odd} k^2\Lambda_k^4 - \sum_{k-even} k^2\Lambda_k^4\right]}} \gg 1$$

The is the second way in which $f_e \gg f$ can be achieved in the current set up.

Supergravity scalar potential [8]

$$V = e^K [D_i W K^{-1}_{ij^*} D_{j^*} W^* - 3|W|^2] + V_D \, ,$$
$$D_i W = W_{,i} + K_{,i} W \, .$$

where K is the Kahler potential.

- We consider here for simplicity a single pair of axions, $\phi_i, \; i = 1, 2$ with opposite shift symmetries where

$$\phi_i = (\rho_i + i a_i)/\sqrt{2}, \; i = 1, 2 \, ,$$

- The following form for the Kähler potential avoids the η problem

$$K = \sum_i \frac{1}{2} (\phi_i + \phi_i^\dagger)^2 \, ,$$

[8] A. H. Chamseddine, R. L. Arnowitt and P. N., Phys. Rev. Lett. **49**, 970 (1982); E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, Nucl. Phys. B **212**, 413 (1983).

Saxion stabilization in SUGRA

- Saxions can be stabilized by imposition of spontaneous symmetry breaking conditions

$$D_i W = 0, \quad i = 1, 2$$

- Choose a new basis

$$b_\pm = \frac{1}{\sqrt{2}}(a_1 \pm a_2)\,.$$

where b_+ is invariant under the shift symmetry while b_- is sensitive to the shift symmetry and is the inflaton

$$W_{sb} = \sum_{n=1}^{q} B_n \left(e^{i\gamma_n \frac{b_-}{\sqrt{2}f}} + e^{-i\gamma_n \frac{b_-}{\sqrt{2}f}}\right).$$

Inflaton potential in supergravity

- Analysis for a single pair of axions.

$$
\begin{aligned}
V_{\text{slow}}(b_-) = \sum_{n=1}^{q}\sum_{m=1}^{q} c_{nm}&\\
&\left(1 - 2\cos(\frac{\gamma_n b_-}{\sqrt{2}f}) + \cos((\gamma_n - \gamma_m)\frac{b_-}{\sqrt{2}f})\right)\\
&+ \sum_{n=1}^{q}\sum_{m=1}^{q} c'_{nm}\\
&\Big[1 - \cos(\gamma_n b_-/\sqrt{2}f) - \cos(\gamma_m b_-/\sqrt{2}f)\\
&+ \frac{1}{2}\cos((\gamma_n - \gamma_m)b_-/\sqrt{2}f) + \frac{1}{2}\cos((\gamma_n + \gamma_m)b_-/\sqrt{2}f)\Big] .
\end{aligned}
$$

Consistency with Planck data [9]

- Some of the quantities of interest for experimental test are the ratio of the tensor and curvature power spectrum and the spectral indices

$$r = \frac{\mathcal{P}_t(k_0)}{\mathcal{P}_R(k_0)} ,$$
$$\mathcal{P}_R(k) = \mathcal{P}_R(k_0)(\frac{k}{k_0})^{n_s(k)-1} ,$$
$$\mathcal{P}_t(k) = \mathcal{P}_t(k_0)(\frac{k}{k_0})^{n_t(k)} .$$

- The current experimental limits from Planck experiment at $k_0 = 0.05\,\mathrm{Mpc}^{-1}$

$$n_s = 0.9645 \pm 0.0049\,(68\%\mathrm{CL})$$
$$r < 0.07\,(95\%\mathrm{CL})$$

while $n_t(k_0)$ is currently not constrained.

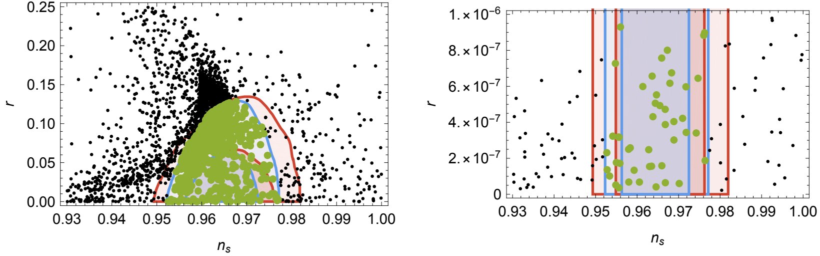
- Require the number of e-foldings

$$N_e = [50, 60].$$

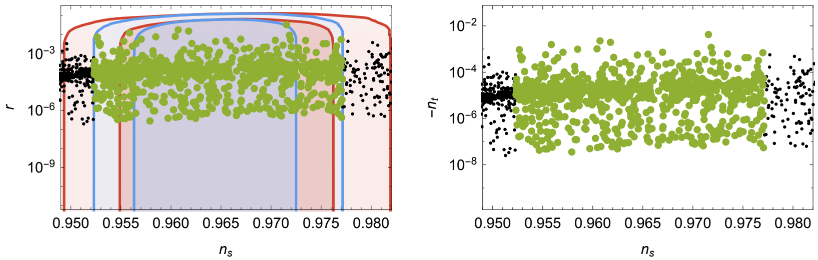
For a mode k, horizon exit occurs at time when $k = RH$ and the Hubble radius is $(RH)^{-1}$.

[9] R. Adam *et al.* [Planck Collaboration], Astron. Astrophys. **594**, A1 (2016); [arXiv:1502.01582 [astro-ph.CO]].
P. A. R. Ade *et al.* [Planck Collaboration], Astron. Astrophys. **594**, A20 (2016) [arXiv:1502.02114 [astro-ph.CO]].

SUSY model



- Left panel:
  - Green and black scatter points: $N_e = [50, 60]$.
  - Green scatter points satisfy experimental bounds on the spectral index n_s and the ratio r.
- The decay constant $f < M_{Pl}$ for all points on the right panel.



Left panel: Display of the ratio r vs the spectral n_s for the supergravity model. Right panel: Display n_t vs n_s for the same parameter space as for the left panel. The green and black parameter points have the same meaning as for the SUSY model.

Non-Gaussianity

- Significant information resides in higher point correlation functions[10].
- For the three point correlation functions the dominant non-Gaussianity arises from the correlation function of three scalar perturbations $\zeta(\vec{k})$ defined by

$$< \zeta(\vec{k}_1)\zeta(\vec{k}_2)\zeta(\vec{k}_3) >= (2\pi)^7\delta^3(\vec{k}_1+\vec{k}_2+\vec{k}_3)\frac{\sum_i k_i^3}{\prod_i k_i^3}\left[-\frac{3}{10}f_{NL}(P_k^\zeta)^2\right]$$

- For equilateral triangle $k_1 = k_2 = k_3$,

$$f_{NL} \sim c(\frac{1}{c_s^2} - 1) + \cdots , c = O(1).$$

- c_s is sound speed

$$c_s^2 = \frac{dp}{d\rho}.$$

For models with canonical kinetic energy, $f_{NL} = 0$ and one needs $c_s \ll 1$ for significant non-Gaussianity to arise. To generate significant non-Gaussianity one needs models non-canonical kinetic terms. One possibility is DBI.

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[10] J. M. Maldacena, JHEP **0305**, 013 (2003); M. Alishahiha et al., Phys. Rev. D **70**, 123505 (2004); D. Seery and J. E. Lidsey, JCAP **0506**, 003 (2005); D. H. Lyth and Y. Rodriguez, Phys. Rev. Lett. **95**, 121302 (2005); X. Chen et al., JCAP **0701**, 002 (2007).

Example of a simple DBI Lagrangian is

$$L_{DBI} = -\frac{1}{f}\sqrt{1 + f\partial_\mu\phi\partial_\nu\phi} - V(\phi)$$

We consider here the SUSY case for a two field model[11]

$$\mathcal{L} = \mathcal{L}_D + \mathcal{L}_F,$$

$$\begin{aligned}
\mathcal{L}_D = & \int d^4\theta \left(\Phi_1\Phi_1^\dagger + \Phi_2\Phi_2^\dagger\right) , \\
& + \int d^4\theta \frac{\alpha_1}{16T} \left(D^\alpha\Phi_1 D_\alpha\Phi_1\right)\left(\overline{D}^{\dot{\alpha}}\Phi_1^\dagger \overline{D}_{\dot{\alpha}}\Phi_1^\dagger\right) G(\phi) \\
& + \int d^4\theta \frac{\alpha_1}{16T} \left(D^\alpha\Phi_2 D_\alpha\Phi_2\right)\left(\overline{D}^{\dot{\alpha}}\Phi_2^\dagger \overline{D}_{\dot{\alpha}}\Phi_2^\dagger\right) G(\phi) \\
& + \int d^4\theta \frac{\alpha_2}{16T} \left(D^\alpha\Phi_1 D_\alpha\Phi_1\right)\left(\overline{D}^{\dot{\alpha}}\Phi_2^\dagger \overline{D}_{\dot{\alpha}}\Phi_2^\dagger\right) G(\phi) \\
& + \int d^4\theta \frac{\alpha_2}{16T} \left(D^\alpha\Phi_2 D_\alpha\Phi_2\right)\left(\overline{D}^{\dot{\alpha}}\Phi_1^\dagger \overline{D}_{\dot{\alpha}}\Phi_1^\dagger\right) G(\phi) \\
& + \int d^4\theta \frac{\alpha_3}{16T} \left(D^\alpha\Phi_1 D_\alpha\Phi_2\right)\left(\overline{D}^{\dot{\alpha}}\Phi_1^\dagger \overline{D}_{\dot{\alpha}}\Phi_2^\dagger\right) G(\phi)
\end{aligned} \tag{2}$$

11 PN, Piskunov, in preparation. A semi-quantitative analysis for a single field DBI model has been considered before by S. Sasaki etal, Phys. Lett. B **718**, 1 (2012).

Function $G(\phi)$

$$G(\phi) = \frac{1}{T} \frac{1}{1 + A + \sqrt{(1+A)^2 - B}}$$

T is a parameter of the dimension of (mass)^4 . A and B are defined by

$$\begin{aligned}
A &= (\partial_a \phi_1 \partial^a \phi_1^* + \partial_a \phi_2 \partial^a \phi_2^*)/T \\
B &= \alpha_1 (\partial_a \phi_1 \partial^a \phi_1 \partial_b \phi_1^* \partial^b \phi_1^* + \partial_a \phi_2 \partial^a \phi_2 \partial_b \phi_2^* \partial^b \phi_2^*)/T^2 \\
&\quad + \alpha_2 (\alpha_3 \partial_a \phi_1 \partial^a \phi_1 \partial_b \phi_2^* \partial^b \phi_2^* + \partial_a \phi_2 \partial^a \phi_2 \partial_b \phi_1^* \partial^b \phi_1^*)/T^2 \\
&\quad + \alpha_3 (\partial_a \phi_1 \partial^a \phi_2 \partial_b \phi_1^* \partial^b \phi_2^*)/T^2
\end{aligned}$$

We will illustrate here by considering the case keeping just the α_1 terms. In this case the equations for the F-terms are cubic which can be solved analytically.

$$F_k^3 + p_k F_k + q_k = 0\,, k = 1, 2\,,$$

where p_k, q_k are defined by

$$p_k = \left(\frac{\partial W}{\partial \varphi_k}\right)^{-1} \frac{\partial W}{\partial \varphi_k} \frac{1 - 2G(\varphi)\,\partial_\mu \varphi_k \partial^\mu \varphi_k}{2G(\varphi)}$$

$$q_k = \frac{1}{2G(\varphi)} \left(\frac{\partial W}{\partial \varphi_k}\right)^{-1} \left(\frac{\partial W}{\partial \varphi_k}\right)^2$$

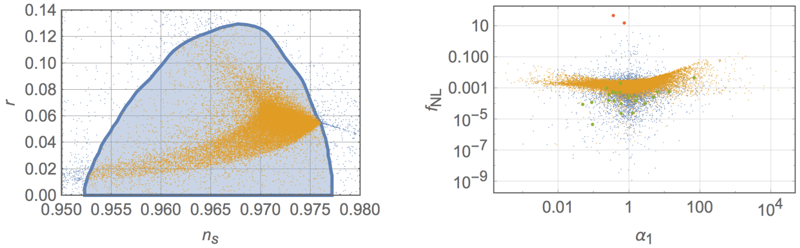
Since F_k satisfies a cubic equation, there are three roots which are given by

$$\begin{aligned} F_k =& \omega^j \left(-\frac{q_k}{2} + \sqrt{\left(\frac{q_k}{2}\right)^2 + \left(\frac{p_k}{3}\right)^3}\right)^{1/3} \\ &+ \omega^{3-j} \left(-\frac{q_k}{2} - \sqrt{\left(\frac{q_k}{2}\right)^2 + \left(\frac{p_k}{3}\right)^3}\right)^{1/3}. \end{aligned}$$

where ω is the cube root of unity and $j = 0, 1, 2$. Physical root is $j = 0$.

SUSY DBI[12]

$$\mathcal{L} = T - T\sqrt{(1+A)^2 - B}$$
$$+ \sum_{i=1}^{2}\Big[F_i F_i^* + \alpha G(\phi)\Big[(-2F_iF_i^*\partial_a\phi_i\partial^a\phi_i^* + F_i^2F_i^{*2})\Big)\Big]$$
$$+ \Big(\frac{\partial W}{\partial\varphi_i}F_i + h.c.\Big)\Big]$$



r can approach the experimental upper limit for SUSY DBI models.

12 PN, M. Piskunov, in preparation

Conclusion

- We have proposed a new mechanism for the enhancement of the axion decay constant within SUSY/SUGRA models with a $U(1)$ symmetry broken by instanton type effects.
- A class of SUSY/supergravity axion models are proposed consistent with data with a sub-Planckian decay constant.
- The mechanism can also be extended to include non-canonical kinetic energy terms via Dirac-Born-Infeld. For SUSY DBI one can get r to be as large as the upper limit from experiment will allow.
- However, getting non-Gaussianity to be in the observable range requires further exploration.

Stronger case for SUSY post Higgs boson discovery at $m_h \sim 125$ GeV

- The measurement of the Higgs boson at 125 GeV gives further support for supersymmetry. This is because of vacuum stability. For large field configurations where $h >> v$ the Higgs potential is governed by the quartic term
$$V_h \sim \lambda_{eff} h^4.$$
Vacuum stability depends critically on the top mass. A larger top mass makes the vacuum more unstable.

An advanced precision analysis[13] including two-loop matching, three-loop renormalization group evolution, and pure QCD corrections through four loops gives an upper bound on the top pole mass for SM stability up to the Planck mass scale of

$$\boldsymbol{m_t^{\mathrm{cri}} = (171.54 \pm 0.30^{+0.26}_{-0.41})\ \mathrm{GeV}.}$$

$$\boldsymbol{m_t^{\mathrm{exp}} = (173.21 \pm 0.51 \pm 0.71)\ \mathrm{GeV}}$$

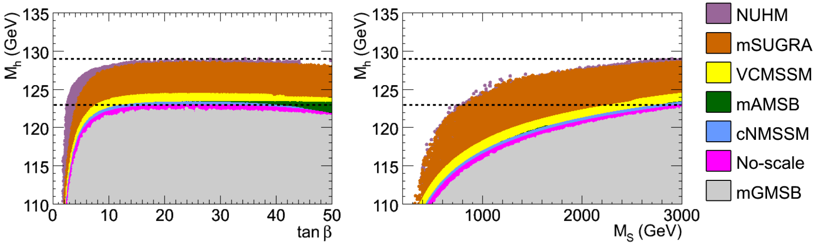
This makes the vacuum stable only up to about $\mathbf{10^{10} - 10^{11}}$ GeV.

In models based on supersymmetry with a Higgs mass of $\mathbf{125}$ GeV, the vacuum can be stable up to the Planck scale.

[13] A. V. Bednyakov, B. A. Kniehl, A. F. Pikelner and O. L. Veretin, Phys. Rev. Lett. **115**, no. 20, 201802 (2015) G. Degrassi, S. Di Vita, J. Elias-Miro, J. R. Espinosa, G. F. Giudice, G. Isidori and A. Strumia, JHEP 1208, 098 (2012).

What light can $m_h \sim 125$ GeV shed on the mechanism of SUSY breaking?

A comparison of mSUGRA, mGMSB, mAMSB and others



A. Arbey, M. Battaglia, A. Djouadi and F. Mahmoudi, JHEP 1209 (2012) 107.

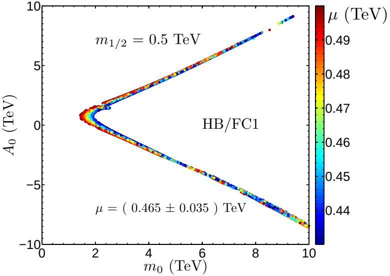
One needs a large A_0 which SUGRA models can generate while other mechanisms have problems with.

Geometry of radiative breaking of the electroweak symmetry

It was noticed in 1998 [14] that the underlying geometry of radiative breaking has an ambiguous sign $\eta = \pm 1$ of the m_0^2 term and thus the geometry can flip from Euclidean to Minkowskian.

$$\mu^2 + \frac{1}{4}M_Z^2 = \eta m_0^2 + A_0^2 + m_{1/2}^2, \quad \eta = \pm 1, 0$$

For $\eta = -1$ one can get large m_0 and A_0 even with low μ. This is just what happens when you explore the soft parameter space.



$\eta = -1$ case: S. Akula, M. Liu, P.N., and G. Peim, Phys. Lett. B **709**, 192 (2012).

[14] Chan, Chattopadhyay, PN, PRD D58, 096004 (1998); H. Baer, C. Balazs, A. Belyaev, T. Krupovnickas and X. Tata, JHEP **0306**, 054 (2003); J. L. Feng, K. T. Matchev and T. Moroi, Phys. Rev. Lett. **84**, 2322 (2000); D. Feldman, G. Kane, E. Kuflik and R. Lu, Phys. Lett. B **704**, 56 (2011).

Conclusion

- Currently there is no paradigm aside from SUSY/Supergravity unification that can extrapolate physics from the electroweak scale to the grand unification scale consistent with the existing data.

- The case for SUSY is stronger now than before the discovery of the Higgs boson mass at 125 GeV.

Extra slides

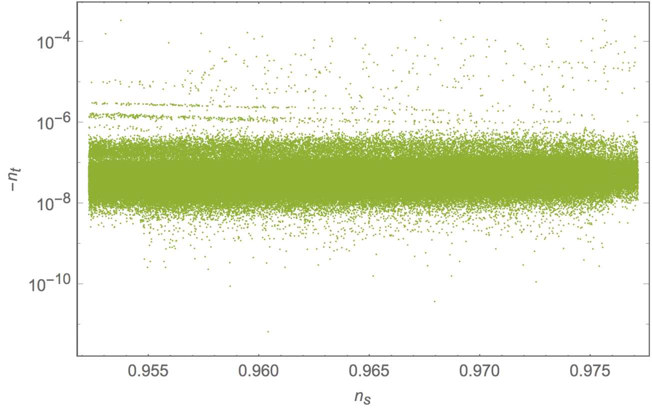


Figure: Monte-Carlo analysis in (n_s, n_t) space for the same data set as in previous figure. Note the log scale on the $-n_t$ axis.

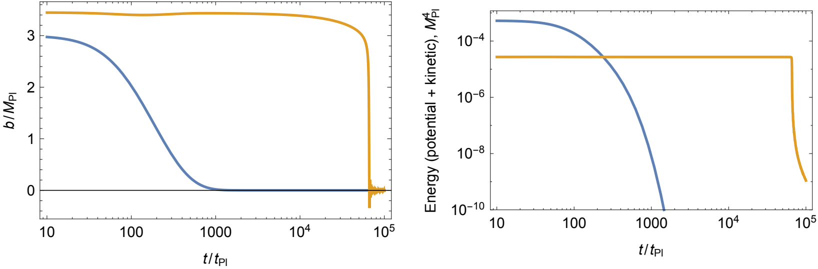


Figure: Fast vs. slow field evolution for the case of Fig.(**??**). Left panel: fast and slow field components as a function of time. Right panel: energy of the slow and fast field components as a function of time. Slow field energy is defined as $E_{\rm slow} = V_{\rm slow}\left(b_-\right) + \frac{1}{2}\dot{b_-}^2$. Fast field energy is defined as $E_{\rm fast} = V_{\rm full}\left(b_+, b_-\right) - V_{\rm slow}\left(b_-\right) + \frac{1}{2}\dot{b_+}^2$. One can see that the fast component starts larger, but then the slow component overtakes it. Note the logarithmic scale of time. Here $A_1 = 10^{-4} M_{Pl}^2$.