July 4, 2019

Axionic Inflation and the Weak Gravity Conjecture

PASCOS2019, Manchester, July 1-5, 2019

PN
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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].

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[1] A. H. Guth, Phys. Rev. D **23**, 347 (1981); A. A. Starobinsky, Phys. Lett. B **91**, 99 (1980)
A. D. Linde, Phys. Lett. **108B**, 389 (1982). A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. **48**, 1220 (1982).

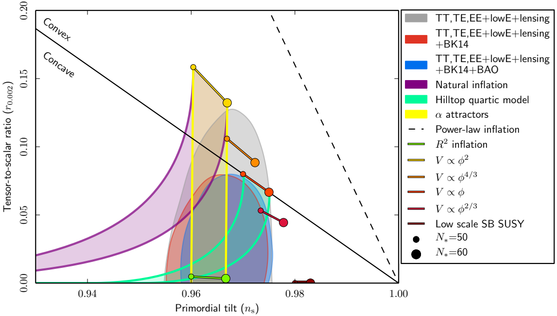


Figure: Limits on the tensor-to-scalar ratio, $r_{0.002}$ as a function of n_S in the ΛCDM model at 95% CL, from Planck alone (grey area) or including BICEP2/Keck data 2014 (red) and BAO (blue). From Y. Akrami *et al.* [Planck Collaboration], arXiv:1807.06205 [astro-ph.CO].

Current limits

$$n_s = 0.9649 \pm 0.0042 \ (68\% \ \mathrm{CL}) \ ,$$

$$r < 0.064 \ (95\% \ \mathrm{CL}) \ ,$$

n_t not constrained.

Axionic inflation

An early inflationary model is the so-called natural inflation where the inflaton is an axion with potential [2]

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

where f is the axion decay constant. For successful inflation one requires

$$f > 5M_{Pl}.$$

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

- The weak gravity conjecture[3] indicates $f < M_{Pl}$.
- Analyses for periodic fields in strings indicate that $f < M_{Pl}$ is the norm (except for some anomalous cases) [4].
- Numerical estimates of the decay constant in strings give the range for f of $(10^{16} - 10^{18})$ GeV[5] .

[3] N. Arkani-Hamed, L. Motl, A. Nicolis and C. Vafa, JHEP **0706**, 060 (2007).

[4] T. Banks, M. Dine, P. J. Fox and E. Gorbatov, JCAP **0306**, 001 (2003).

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

Effective theory of inflation

Inflation occurs at scales far below the Planck scale. Here the effective Lagrangian that governs inflation would be very different from the microscopic Lagrangian.

While $f < M_{Pl}$ may hold for the microscopic Lagrangian, the decay constant that enters in the effective theory could be different which gives the hope of generating an $f_e > M_{Pl}$.

There are several proposals which try to generate such effective theories. Two example are

- Alignment mechanism
- Coherent enhancement mechanism (CEM)

Alignment mechanism [6]

One suggestion to realize $f_e >> f$ 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

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

- A flat direction is generated if

$$\frac{f_1}{f_2} = \frac{f_3}{f_4}.$$

  One then considers deviations from the constraint to lift the flat direction and generate an effective decay constant for the inflaton.
- Difficult to implement in strings with stabilized moduli[7].

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

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

Axion Landscape

- The axions we consider are not QCD axions for which the decay constant lies in the range

$$10^9 \text{ GeV} < f < 10^{12} \text{ GeV}.$$

- Rather our axions are string axions where f typically lies in the range $10^{16} \text{ GeV} - 10^{18} \text{ GeV}$.
- For generality consider a landscape of m number of chiral fields charged under a $U(1)$ shift symmetry and a corresponding number of chiral fields which are oppositely charged.
- There are then $2m$ number of axionic fields

$$a_1, a_2, \cdots a_m; \bar{a}_1, \bar{a}_2, \cdots \bar{a}_m$$

We can then construct the following linear combinations

$$
\begin{aligned}
b_k &= \frac{a_{k+1}}{f_{k+1}} - \frac{a_1}{f_1}, \quad k = 1, 2, \cdots, m-1\,, \\
\overline{b}_k &= \frac{\overline{a}_{k+1}}{\overline{f}_{k+1}} - \frac{\overline{a}_1}{\overline{f}_1}, \quad k = 1, 2, \cdots, m-1\,, \\
b_+ &= \frac{a_1}{f_1} + \frac{\overline{a}_1}{\overline{f}_1}\,, \\
b_- &= \frac{1}{\sqrt{\sum_{k=1}^m f_k^2 + \sum_{k=1}^m \overline{f}_k^2}} \left(\sum_{k=1}^m f_k a_k - \sum_{k=1}^m \overline{f}_k \overline{a}_k \right)\,.
\end{aligned}
\tag{1}
$$

Thus the first three equations in Eq.(1) give us $2(m-1)+1 = 2m-1$ linear combinations of axionic fields which are invariant under the shift symmetry while the last one gives us the combination of axionic fields which is sensitive to shift symmetry. b_- is the pseudo-Nambu-Goldstone-Boson (pNGB).

pNGB is the inflaton[8]

- Consider a superpotential of the form

$$W = W_s + W_{sb}$$

where W_s is invariant under the shift symmetry and W_{sb} breaks the shift symmetry. The stability conditions $W_{,i} = 0$ will generate large masses for all the fields except b_- while b_- receives mass only from W_{sb}.

- The potential of the axions can be decomposed into two parts so that

$$V = V_{\text{fast}}(b_k/f_{eN}, \bar{b}_k/f_{eN}, b_+/f_{eN}) + V_{\text{slow}}(b_-/f_{eN})$$

Inflation is controlled only by b_- and the remaining fields play no role. When $f_k = \bar{f}_k = f$, $N = 2m$,

$$f_{eN} = \sqrt{\sum_{k=1}^{m} f_k^2 + \sum_{k=1}^{m} \bar{f}_k^2} = \sqrt{N} f \,. \quad \text{N-flation}$$

[8] P.N. and M. Piskunov, JHEP **1803**, 121 (2018) [arXiv:1712.01357 [hep-ph]].

General analysis of Coherent Enhancement Mechanism for f_e

Consider an axionic inflaton with a Lagrangian of the form

$$\mathcal{L} = \frac{1}{2}\dot{\phi}^2 - V(\phi)$$

Let us assume that $\phi \sim \phi_0$ is the point where one has horizon exit and simulate the potential near this point by

$$V_e(\phi) = \Lambda^4 \left(1 - \cos(\frac{\phi_0}{f_e})\right)$$

It then follows that

$$f_e = \frac{V(\phi_0)}{\sqrt{V'^2(\phi_0) - 2V(\phi_0)V''(\phi_0)}} \tag{2}$$

Effective axion decay constant f_e in terms of r and n_s

$$\epsilon = \frac{M_{Pl}^2}{2}\left(\frac{V'(\phi_0)}{V(\phi_0)}\right)^2 , \quad \eta = M_{Pl}^2 \frac{V''(\phi_0)}{V(\phi_0)} .$$

This leads to

$$f_e = \frac{M_{Pl}}{\sqrt{2(\epsilon - \eta)}} = \frac{M_{Pl}}{\sqrt{1 - n_s - r/4}} .$$

where we used the relation: $n_s = 1 - 6\epsilon + 2\eta , \quad r = 16\epsilon .$

Planck data:

$$n_s = 0.9649 \pm 0.0042 \, (68\% \, \mathrm{CL}) , \quad r < 0.064 \, (95\% \, \mathrm{CL})$$

gives model-independent bounds on f_e

$$4.9 \leq f_e/M_{Pl} \leq 10.0 \, (95\% \, \mathrm{CL}) . \tag{3}$$

Illustration of Coherent Enhancement Mechanism (CEM) [9]

CEM arises from constructive and destructive interference among several terms in the potential. Consider, for example, the potential

$$V = \sum_{k=1}^{n} \Lambda_k^4 \left(1 - \cos\left(\frac{k\phi}{f}\right)\right) .$$

Assume that near the horizon exit $\frac{\phi}{f} = \pi$, then using Eq. (2)

$$f_e/f = \frac{\sqrt{\sum_{k \in odd} \Lambda_k^4}}{\sqrt{\sum_{k \in odd} k^2 \Lambda_k^4 - \sum_{k \in even} k^2 \Lambda_k^4}} .$$

A cancellation between the odd and even sums in the denominator leads to $f_e/f > 1$.

[9] P.N. and M. Piskunov, arXiv:1906.02764 [hep-ph].

Three classes of models

- SUSY models
- Supergravity models
- Supersymmetric Dirac-Born -Infeld models

SUSY model

Suppose we have a set of fields Φ and $\bar{\Phi}$ which carry opposite charges under a shift symmetry. We parametrize the scalar components so that

$$\phi = (f+\rho)e^{ia/f}, \quad \bar{\phi} = (\bar{f}+\bar{\rho})e^{i\bar{a}/f} \,,$$

This allows us to write a non-trivial superpotential which can stabilize the saxions[10]

$$\begin{aligned} W &= W_s + W_{sb} \\ W_{sb} &= \sum_{l=1}^{q} \left(A_l \Phi^l + \bar{A}_l \bar{\Phi}^l \right) . \end{aligned}$$

W_s is symmetry preserving and W_{sb} breaks the shift symmetry. Saxions are stabilized via constraints

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

[10]Similar to superpotentials in analysis of an ultralight axion:
J. Halverson, C. Long and P.N., Phys. Rev. D **96**, no. 5, 056025 (2017).

Slow roll potential with stabilized saxions

We define linear combinations $b_\pm$ of $a, \bar{a}$ so that b_+ is invariant under the shift symmetry and b_- is not.

$$b_\pm = \frac{1}{\sqrt{2}}(a \pm \bar{a}).$$

b_+ undergoes fast roll and b_- is the infation which undergoes slow roll and controls inflation [11]

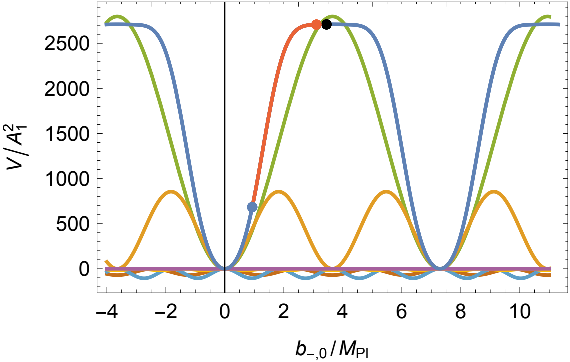
$$V(b) = V_{\text{fast}} + V_{\text{slow}}(b_-)$$

$$V_{\text{slow}} = \sum_{r=1}^{q} C_r \left(1 - \cos\left(\frac{r}{\sqrt{2}f} b_-\right)\right) + \sum_{s=1}^{q} \sum_{r=s+1}^{q} C_{rs}\left(1 - \cos\left(\frac{r-s}{\sqrt{2}f} b_-\right)\right)$$

[11] PN., M. Piskunov, JHEP **1803**, 121 (2018) [arXiv:1712.01357 [hep-ph]].

Emergence of a locally flat potential

Red curve: Inflation potential.



●: Initial value of field.

●: Horizon exit.

●: End of inflation.

Emergence of effective decay constant

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

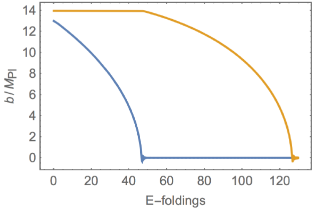
$$
\begin{aligned}
V\left(b_{-}\right) = {} & 1398.96\left(1-\cos\left(\frac{b_{-}}{\sqrt{2}f}\right)\right)+427.466\left(1-\cos\left(\frac{2b_{-}}{\sqrt{2}f}\right)\right) \\
& -35.4939\left(1-\cos\left(\frac{3b_{-}}{\sqrt{2}f}\right)\right)-52.9837\left(1-\cos\left(\frac{4b_{-}}{\sqrt{2}f}\right)\right)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

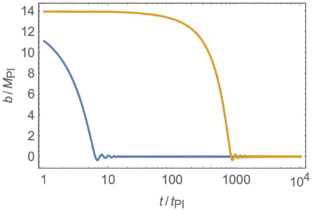
- Superposition gives local flatness.

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

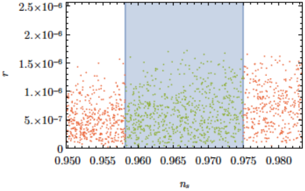
$$
f_e/f = \frac{\sqrt{\sum_{k-odd}\Lambda_k^4}}{\sqrt{\left[\sum_{k-odd}k^2\Lambda_k^4-\sum_{k-even}k^2\Lambda_k^4\right]}} = \frac{\sqrt{1355}}{\sqrt{872-839}} \simeq 6.4
$$

Fast roll: b_+ (blue) Slow roll: b_- (red)

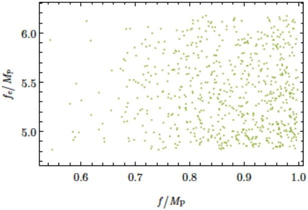




r vs n_s



f_e vs f



Axion inflation in Supergravity

We can extend the analysis to supergravity where the scalar potential has the form [12]

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

- As for the global SUSY case, here too we consider a single pair of chiral fields, $\phi_i, \; i = 1, 2$ 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

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

[12] 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

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

- As in global SUSY $a_- = \frac{1}{\sqrt{2}}(a_1 - a_2)$ is the inflaton and the superpotential that involves the inflaton is

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

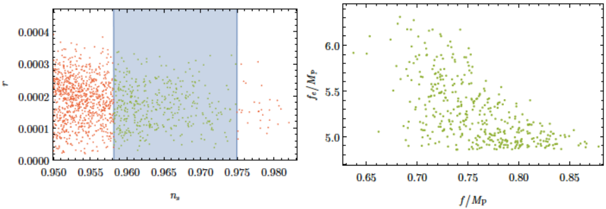
Inflaton potential in supergravity

- Analysis for the case $q = 3$ gives

$$V\left(a_-\right) = M_P^4 e^{2f^2/M_P^2} \sum_{k=1}^{6} C_k \left(1 - \cos\left(\frac{ka_-}{\sqrt{2}f}\right)\right),$$

where C_k are given in terms of A_i (i=1-3). One has a linear combination of six cosine terms and their superposition makes the model very different from the natural inflation model.

r vs n_s $\qquad$ f_e vs f



Supersymmetric Dirac-Born-Infeld models

Example of a simple non-supersymmetric DBI action is

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

Early work on supersymmetric DBI model with single field was done by Rocek (1997), Tseytlin (1999), Sasaki, Yamaguchi, Yokoyama (2012). Here we consider a two field supersymmetric DBI model with fields (Φ_1, Φ_2) model[13]
$\mathcal{L} = \mathcal{L}_D + \mathcal{L}_F$, where

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

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

where T is a parameter of the dimension of (mass)^4 .

$$A = (\partial_a\phi_1\partial^a\phi_1^* + \partial_a\phi_2\partial^a\phi_2^*)/T$$

$$B = \alpha_1(\partial_a\phi_1\partial^a\phi_1\partial_b\phi_1^*\partial^b\phi_1^* + \partial_a\phi_2\partial^a\phi_2\partial_b\phi_2^*\partial^b\phi_2^*)/T^2$$

[13] PN, M. Piskunov, JHEP **1902**, 034 (2019) [arXiv:1807.02549 [hep-ph]].

F-equations

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

where p_k, q_k are defined by

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

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

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

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

SUSY DBI [14]

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

Slow roll parameters for DBI (Maldacena 2002, Seery 2005, Chen et al 2005, Lyth 2005)

$$\epsilon_H = -\frac{\dot{H}}{H^2}, \quad \eta_H = \frac{\dot{\epsilon}_H}{\epsilon_H H}, \quad s_H = \frac{\dot{c}_s}{c_s H}. \tag{4}$$

Further, the spectral indices n_s, n_t in this case are given by

$$n_s = 1 - 2\epsilon_H - \eta_H - s_H ,$$
$$n_t = -2\epsilon_H , \quad r = -8 c_s n_t . \tag{5}$$

C_s is the speed of sound in the medium.

[14] PN, M. Piskunov, arXiv:1906.02764 [hep-ph].

Explicit form of DBI Lagrangian

$$\mathcal{L} = \mathcal{L}_I + \mathcal{L}_{II}\,. \tag{6}$$

$\mathcal{L}_I$ is given by

$$\mathcal{L}_I = T\left(1 - \sqrt{1 - \frac{\dot{a}_-^2}{T} + \frac{(2-\alpha_1)\,\dot{a}_-^4}{8T^2}}\right). \tag{7}$$

$\mathcal{L}_{II}$ is more complicated:

$$\begin{aligned}\mathcal{L}_{II} = T\Bigg(& 2\mathcal{F}_+^2 + 2\mathcal{F}_-^2 - \frac{4}{3\alpha_1}\left(\mathcal{T} + (\alpha_1 - 1)\,\frac{\dot{a}_-^2}{4T}\right) + 4k\left(\mathcal{F}_+ + \mathcal{F}_-\right) \\ &+ \frac{\alpha_1}{\mathcal{T} - \dot{a}_-^2/(4T)}\left(2\left(\mathcal{F}_+^2 + \mathcal{F}_-^2 - \frac{2}{3\alpha_1}\left(\mathcal{T} + (\alpha_1 - 1)\,\frac{\dot{a}_-^2}{4T}\right)\right)\frac{\dot{a}_-^2}{4T} + \mathcal{F}_+^4 + \mathcal{F}_-^4 \right. \\ &\left. + \frac{2}{3\alpha_1^2}\left(\mathcal{T} + (\alpha_1 - 1)\,\frac{\dot{a}_-^2}{4T}\right)^2 - \frac{4}{3\alpha_1}\left(\mathcal{T} + (\alpha_1 - 1)\,\frac{\dot{a}_-^2}{4T}\right)\left(\mathcal{F}_+^2 + \mathcal{F}_-^2\right)\right)\Bigg),\end{aligned} \tag{8}$$

where

$$\mathcal{T} = \frac{1}{2}\left(1 + \sqrt{1 - \frac{\dot{a}_-^2}{T} + \frac{(2-\alpha_1)\,\dot{a}_-^4}{8T^2}}\right), \tag{9}$$

$$k = \tilde{\beta}\sqrt{\sum_{m,n} mn\mathcal{G}_m\mathcal{G}_n\left(1 - \cos\left(\frac{a_- m}{\sqrt{2}f}\right) - \cos\left(\frac{a_- n}{\sqrt{2}f}\right) + \cos\left(\frac{a_-\,(m-n)}{\sqrt{2}f}\right)\right)}, \tag{10}$$

$$\mathcal{F}_\pm = \pm\left(\mp\frac{1}{2\alpha_1}k\left(\mathcal{T} - \frac{\dot{a}_-^2}{4T}\right) + \sqrt{\frac{1}{4\alpha_1^2}k^2\left(\mathcal{T} - \frac{\dot{a}_-^2}{4T}\right)^2 + \frac{1}{27\alpha_1^3}\left(\mathcal{T} + (\alpha_1 - 1)\,\frac{\dot{a}_-^2}{4T}\right)^3}\right)^{1/3}, \tag{11}$$

$$\mathcal{G}_k = \frac{A_k 2^{1/2(1-k)}}{\tilde{\beta}\sqrt{T}f^{1-k}}\,. \tag{12}$$

Here $\tilde{\beta}$ is an arbitrary dimensionless parameter which we choose such that $\mathcal{G}_k \sim 1$, and which determines the scale of symmetry breaking terms relative to T.

Friedman equations for DBI

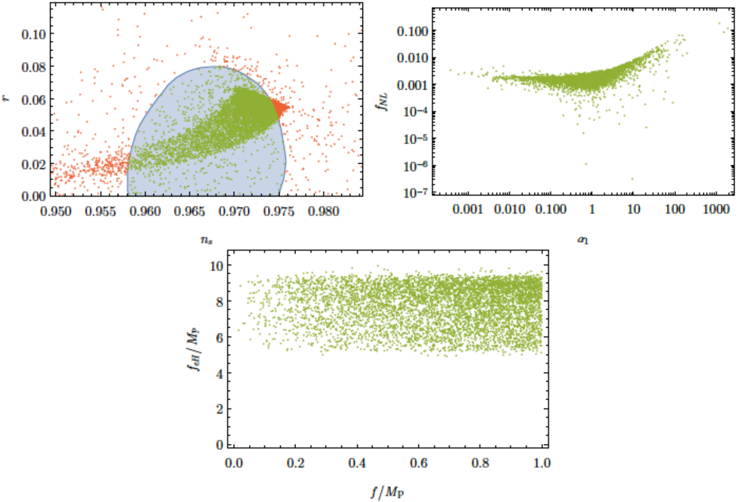
$$3M_{\mathrm{P}}^2\frac{\dot{R}^2}{R^2} = 2\dot{a}_-^2\frac{\partial L}{\partial \dot{a}_-^2} - L.$$

$$2\left[2\frac{\partial^2 L}{\partial \dot{a}_-^2\partial \dot{a}_-^2}\dot{a}_-^2 + \frac{\partial L}{\partial \dot{a}_-^2}\right]\ddot{a}_- - \frac{\partial L}{\partial a_-} + 6H\dot{a}_-\frac{\partial L}{\partial \dot{a}_-^2} + 2\frac{\partial^2 L}{\partial a_-\partial \dot{a}_-^2}\dot{a}_-^2 = 0$$

The potential does not enter by itself, and it is the Lagrangian that governs the Friedman equations.

Approximations one uses in flat potential models are not possible here. Analysis of DBI requires numerical simulations.

Result of numerical simulations[15]



[15] PN, M. Piskunov, JHEP **1902**, 034 (2019) [arXiv:1807.02549 [hep-ph]].

The theoretical upper limit on r in single field inflation for flat potentials

For flat potentials slow roll gives

$$N = -\frac{1}{M_{Pl}^2}\int \frac{V(\phi)}{V'(\phi)}d\phi \sim -\frac{\Delta\phi}{M_{Pl}\sqrt{2\epsilon}}$$

$$r \sim \frac{8(\Delta\phi)^2}{M_{Pl}^2 N^2}$$

For $\Delta\phi/M_{Pl} < 1$ and $N = [50 - 60]$

$$r < 0.003\,. \tag{13}$$

Thus we find that the theoretical upper limit on r in single field inflation for flat potentials is $O(10^{-3})$.

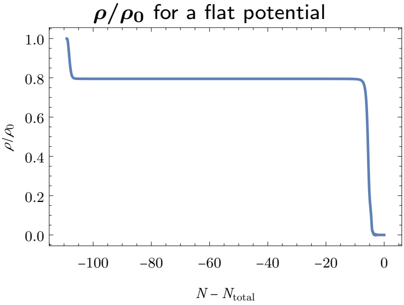
r in DBI

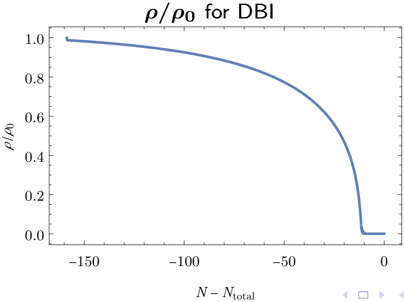
For DBI the density governs the slow-roll and not the potential, i.e.,

$$r = -\frac{8}{\rho}\frac{d\rho}{dN}\,.$$

r can be as large as the current experimental limit of $r = 0.064$ if the slope $d\rho/dN$ is large.

The slope $d\rho/dN$ determines the size of r.

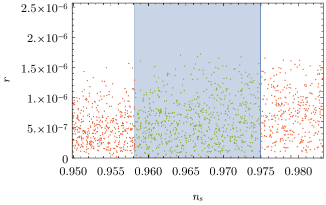




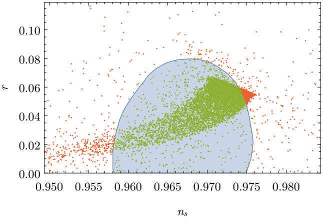
Summary

• Axion inflation with $N = [50 - 60]$ and (n_s, r) in the blue region of Planck data can be achieved with <u>sub-Planckian</u> decay constant consistent with WGC.

SUSY-flat potential



SUSY- DBI



• DBI allows for much larger values of r than the ones in flat potential for single field inflation.

Extra Slides

H at horizon exit

Inflaton mass

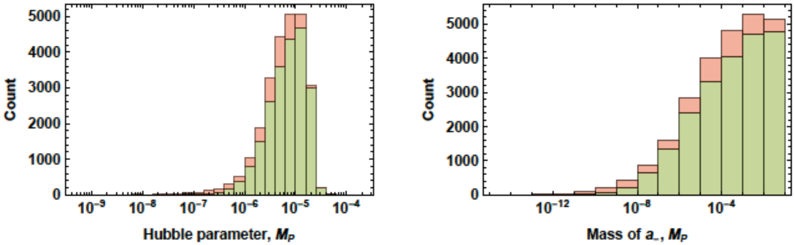


Figure 5. Left panel: a histogram of the values of Hubble parameter at horizon exit for the same data set as in figure 1 for points that are consistent (green) and not consistent (red) with experimental constraints on r and n_s. Right panel: the same as the left panel except the histogram of the inflaton a_- mass is shown instead.