Axion Monodromy Inflation and Related Phenomenology

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arXiv:0907.2916
w/
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arXiv:1002.0833
w/
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arXiv:1106.3335
w/
Mustafa Amin, Richard Easther, Hal Finkel, Mark Hertzberg

ACP Seminar, IPMU, February 15, 2012

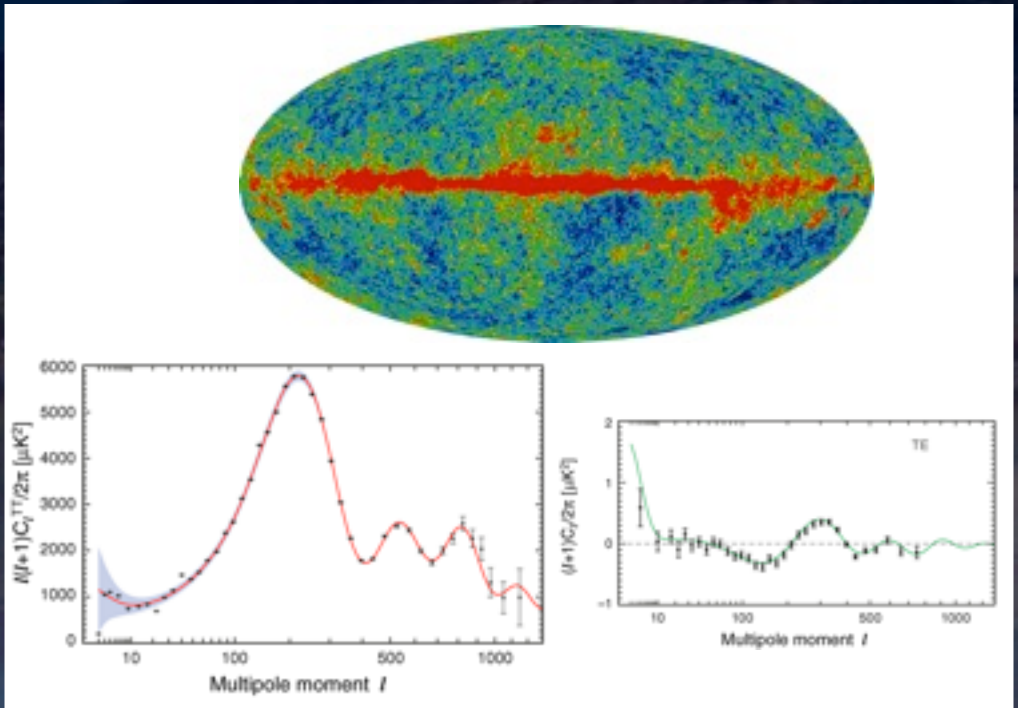
Outline

- Motivation
- Basic Ingredients of Axion Monodromy Inflation
- “Model Independent” Signatures
- “Model Dependent” Signatures
- Conclusions

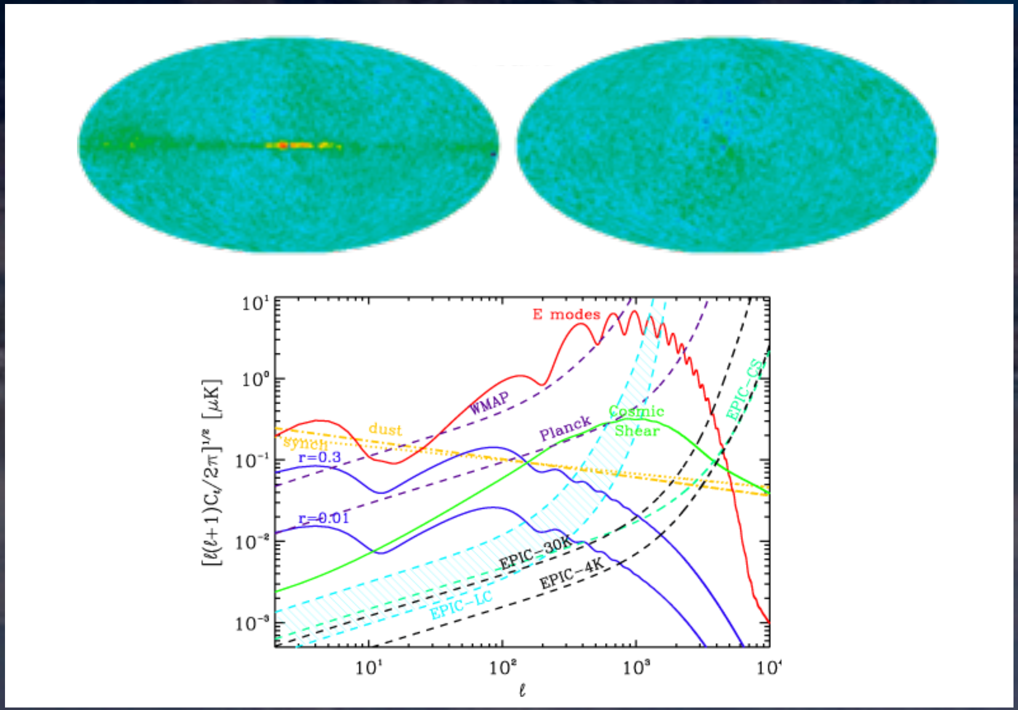
Motivation

(Jarosik et al. 2010)

(Jarosik et al. 2010)



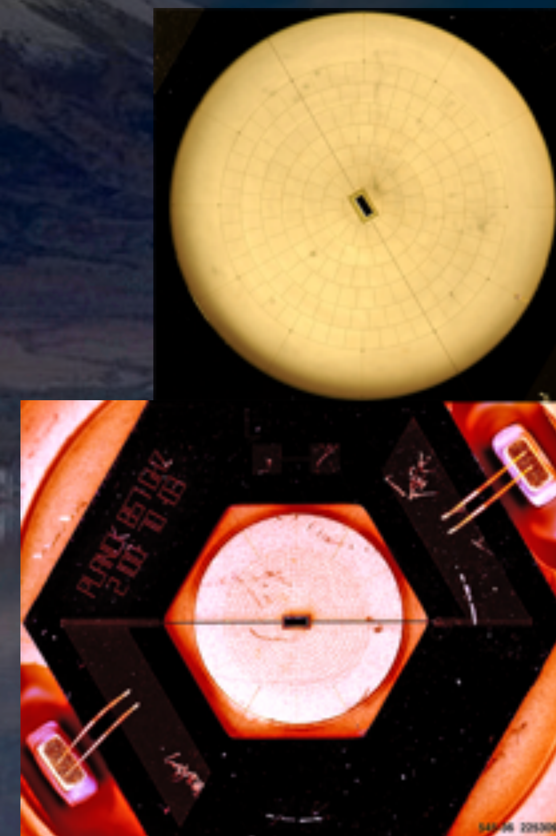
(Larson et al. 2010)



(Bock et al. 2009)

CMB Observables

(Holmes et al. 2008)



To a good first approximation, the detectors are devices that count the number of photons hitting them per unit time per unit area. Some of them are sensitive to polarization.

They then measure the intensity

$$\mathcal{I}(\mathbf{x}, \hat{n}, t, \gamma)$$

CMB Observables

To a good first approximation, the detectors are devices that count the number of photons hitting them per unit time per unit area. Some of them are sensitive to polarization.

(Holmes et al. 2008)

They then measure the intensity

$$\mathcal{I}(0,\hat{n},t_0,\gamma) = \mathcal{I}_0\left(1+\frac{4\Delta T(\hat{n})}{T_0}\right.$$

$$\left.+\frac{4\sqrt{Q^2(\hat{n})+U^2(\hat{n})}}{T_0}\cos\left(2\gamma(\hat{n})-\arctan\frac{U(\hat{n})}{Q(\hat{n})}\right)\right)$$

CMB Observables

It turns out to be more convenient to trade the temperature T and Stokes parameters Q and U in for multipole coefficients

$$a_{T,\ell m} = \int d^2\hat{n}\, Y_\ell^{m*}(\hat{n}) \Delta T(\hat{n})$$

$$a_{P,\ell m} = \int d^2\hat{n}\, {}_2Y_\ell^{m*}(\hat{n}) \left(Q(\hat{n}) + iU(\hat{n})\right)$$

$$a_{E,\ell\, m} \equiv -(a_{P,\ell\, m} + a^*_{P,\ell\, -m})/2$$

$$a_{B,\ell\, m} \equiv i(a_{P,\ell\, m} - a^*_{P,\ell\, -m})/2$$

CMB Observables

and measure the corresponding angular power spectra

$$C^{\rm obs}_{TT,\ell} \equiv \frac{1}{2\ell+1}\sum_m \left|a^{\rm obs}_{T,\ell\, m}\right|^2$$

$$C^{\rm obs}_{TE,\ell} \equiv \frac{1}{2\ell+1}\sum_m a^{\rm obs}_{T,\ell m}\, a^{\rm obs}_{E,\ell m}{}^{*}$$

$$C^{\rm obs}_{EE,\ell} \equiv \frac{1}{2\ell+1}\sum_m \left|a^{\rm obs}_{E,\ell m}\right|^2$$

$$C^{\rm obs}_{BB,\ell} \equiv \frac{1}{2\ell+1}\sum_m \left|a^{\rm obs}_{B,\ell m}\right|^2$$

CMB Observables

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$$
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\begin{aligned}
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\end{aligned}
\right\} \sim \Delta^2_{\mathcal{R}}(k) = \frac{H^2(t_k)}{8\pi^2\epsilon(t_k)}
$$

scalar and tensor perturbations

$$
C^{\rm obs}_{BB,\ell} \equiv \frac{1}{2\ell+1}\sum_m \left|a^{\rm obs}_{B,\ell m}\right|^2
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CMB Observables

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scalar and tensor perturbations

$$C^{\text{obs}}_{BB,\ell} \equiv \frac{1}{2\ell+1}\sum_m \left|a^{\text{obs}}_{B,\ell m}\right|^2 \Big\} \sim \Delta^2_h(k) = \frac{2H^2(t_k)}{\pi^2}$$

only tensor perturbations

CMB Observables

and measure the corresponding angular power spectra

$$
\left.
\begin{aligned}
C^{\rm obs}_{TT,\ell} &\equiv \frac{1}{2\ell+1}\sum_m \left|a^{\rm obs}_{T,\ell\, m}\right|^2 \\
C^{\rm obs}_{TE,\ell} &\equiv \frac{1}{2\ell+1}\sum_m a^{\rm obs}_{T,\ell m}{a^{\rm obs}_{E,\ell m}}^* \\
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\end{aligned}
\right\} \sim \Delta^2_{\mathcal{R}}(k) = \frac{H^2(t_k)}{8\pi^2\epsilon(t_k)}
$$

$$
r = \frac{\Delta^2_h}{\Delta^2_{\mathcal{R}}}
$$

$$
C^{\rm obs}_{BB,\ell} \equiv \frac{1}{2\ell+1}\sum_m \left|a^{\rm obs}_{B,\ell m}\right|^2 \Bigg\} \sim \Delta^2_h(k) = \frac{2H^2(t_k)}{\pi^2}
$$

Motivation

In single field slow-roll inflation, the energy scale of inflation and the distance traveled by the inflaton in field space are related to the tensor-to-scalar ratio

$$V_{\rm inf}^{1/4} = 1.06 \times 10^{16}\, GeV \left(\frac{r}{0.01}\right)^{1/4}$$

$$\Delta\phi \approx \Delta N \sqrt{\frac{r}{8}} \approx \sqrt{\frac{r}{0.01}}\; M_P$$

If a tensor signal is seen, the inflaton must have moved over a super-Planckian distance in field space* **(Lyth 1996)**

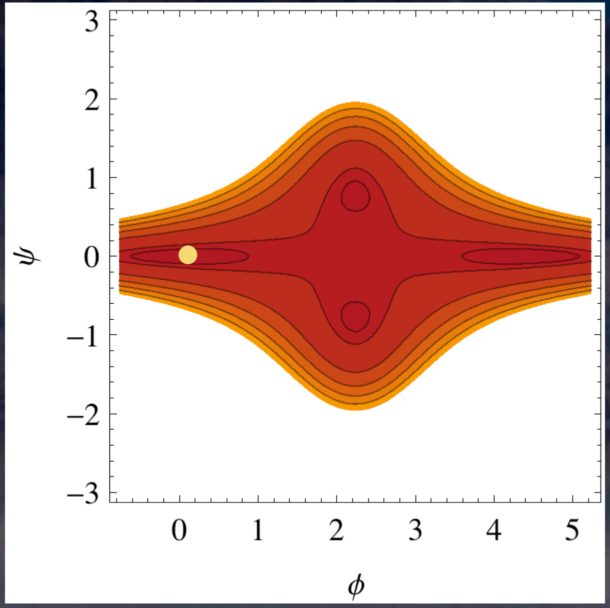
*This is for single field slow-roll models with canonical kinetic term. For a very nice treatment for more general single field models, see arXiv:1111.3040 (Green & Baumann)

Motivation

Motion of the scalar field over super-Planckian distance is hard to control in an effective field theory

$$V(\phi) = V_0 + \frac{1}{2}m^2\phi^2 + \frac{1}{3}\mu\phi^3 + \frac{1}{4}\lambda\phi^4 + \phi^4 \sum_{n=1}^{\infty} c_n \left(\phi/\Lambda\right)^n$$

$$(\Lambda < M_p)$$



The c_n are typically unknown.
Even if they were known, the effective theory is generically expected to break down for $\phi > \Lambda$, e.g. because other degrees of freedom become light.

Motivation

A possible solution:

Use a field with a shift symmetry.
Break the shift symmetry in a controlled way.

The inflaton as an axion

Freese, Frieman, Olinto, PRL 65 (1990)

$$V(\phi) = \Lambda^4 \left[1 + \cos\left(\frac{\phi}{f}\right)\right] \quad \text{with} \quad f \gtrsim M_p$$

However, such large f seem hard to realize string theory.

Banks, Dine, Fox, Gorbatov hep-th/0303252

Motivation

First example of large field inflation in string theory

Silverstein, Westphal, arXiv:0803.3085

and later

McAllister, Silverstein, Westphal, arXiv:0808.0706
Flauger, McAllister, Pajer, Westphal, Xu, arXiv:0907.2916

Berg, Pajer, Sjors, arXiv:0912.1341

For interesting related studies in EFT, see

Kaloper, Lawrence, Sorbo, arXiv:1101.0026

Dubovsky, Lawrence, Roberts, arXiv:1105.3740

Basic Ingredients for Axion Monodromy Inflation

Axion Monodromy Inflation

Origin of the inflaton

In string theory axions arise from integrating gauge potentials over non-trivial cycles in the compactification manifold.

$$b_I(x) = \int_{\Sigma_I^{(2)}} B$$

$$c_\alpha(x) = \int_{\Sigma_\alpha^{(p)}} C^{(p)}$$

In the limit of zero momentum, the couplings of these fields vanish to all orders in string perturbation theory.

Axion Monodromy Inflation

Origin of the potential

Breaking by branes

For definiteness consider a D5-brane wrapping a two-cycle $\Sigma^{(2)}$ of size L$\sqrt{\alpha'}$.

$$S_{\rm DBI} = -\frac{1}{(2\pi)^5 {\alpha'}^3 g_s} \int d^6\xi \sqrt{det(-\varphi^*(G+B))}$$

$$\supset -\frac{\epsilon}{(2\pi)^5 {\alpha'}^2 g_s} \int d^4x \sqrt{{}^{(4)}g} \sqrt{L^4 + b^2}$$

Axion Monodromy Inflation

So the axion has the following potential

$$V(b) = \frac{\epsilon}{(2\pi)^5 {\alpha'}^2 g_s} \sqrt{L^4 + b^2}$$

For large field values in terms of the canonically normalized fields

$$V(\phi) \approx \mu^3 \phi$$

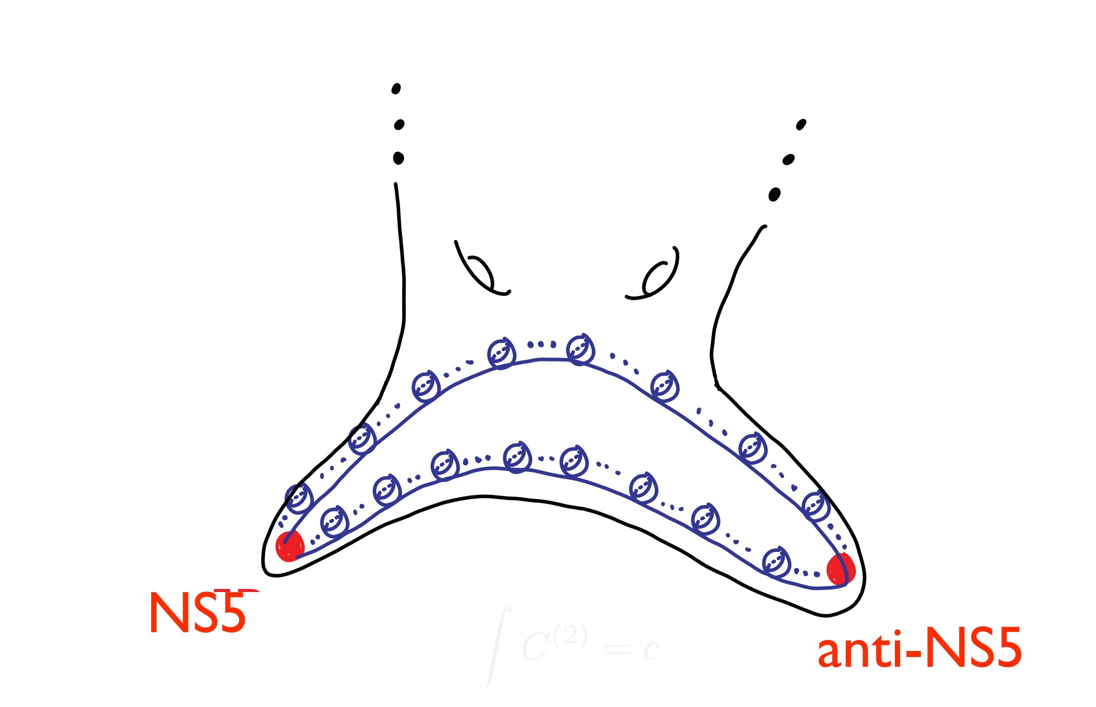
with $$\mu = \frac{\epsilon^{1/3} (2\pi)^3 g_s}{L^{10/3}} M_p$$

Similarly for the $C^{(2)}$ axion in the presence of NS5 branes

Axion Monodromy Inflation

The basic setup

- Type IIB orientifolds with O3/O7
- Stabilize the moduli a la KKLT



Axion Monodromy Inflation

Consistency checks

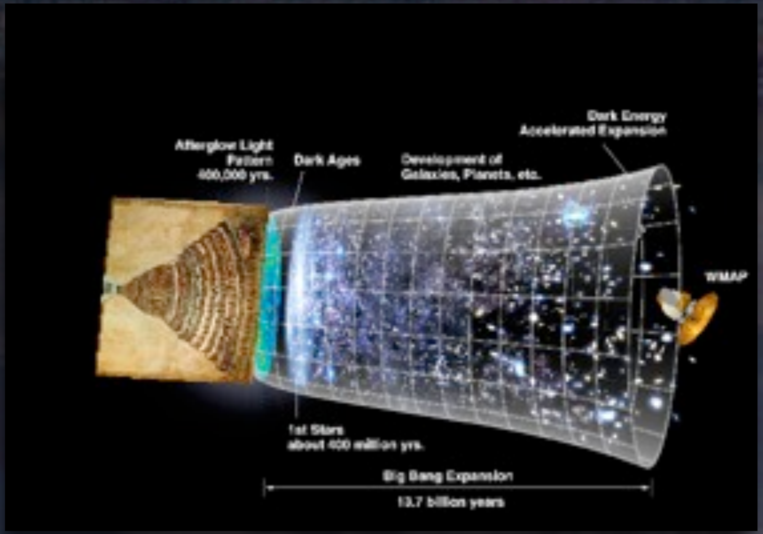
The inflaton potential must be smaller than the potential barriers stabilizing the moduli.

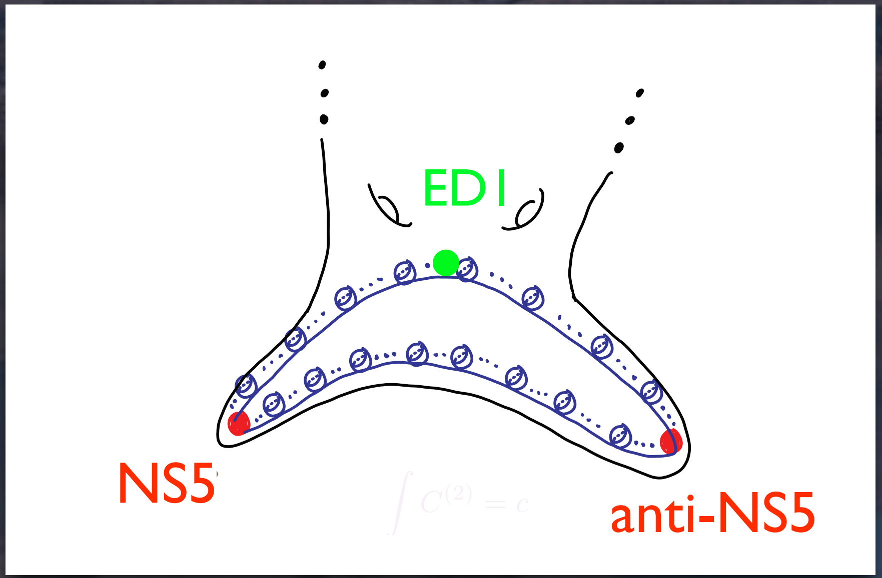
The backreaction on the geometry must be controlled.

Higher derivative corrections must be negligible.

Instanton corrections must be controlled.

Phenomenology





Signatures in the CMB

The low energy effective field theory for Axion Monodromy Inflation is that of a single scalar field with canonical kinetic term, minimally coupled to gravity, with potential

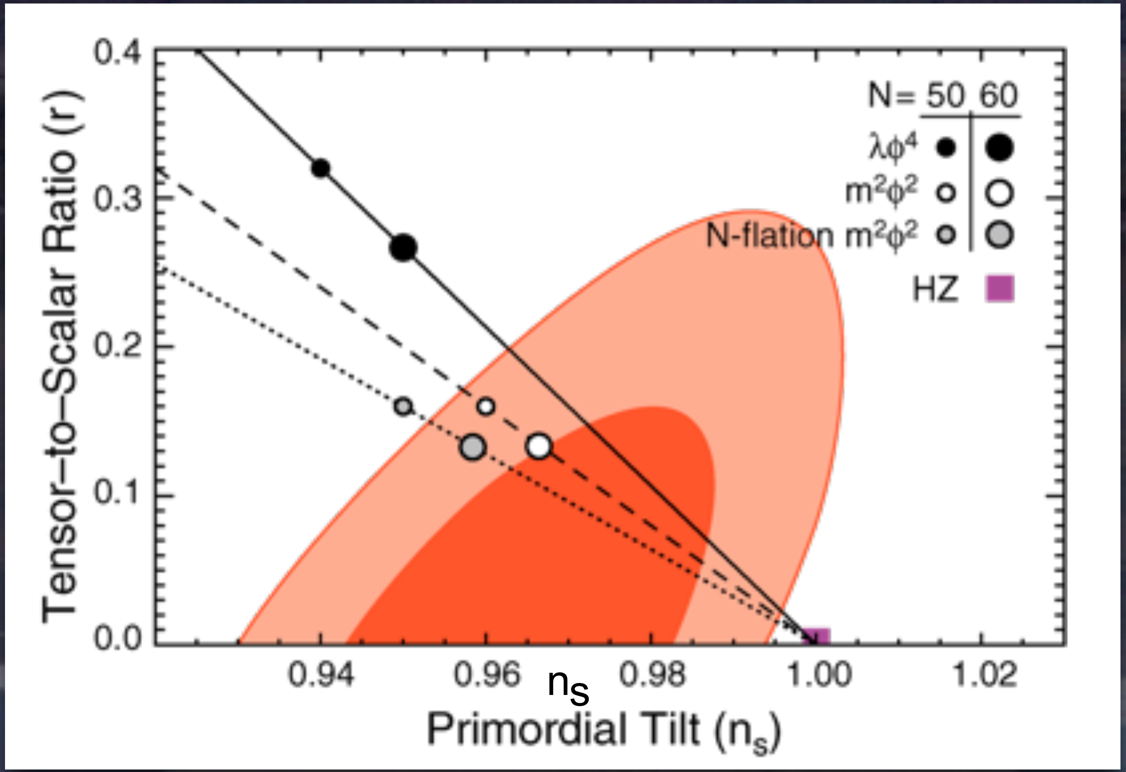
$$V(\phi) = \mu^3\phi + b\mu^3 f\cos(\phi/f)$$

possibly with additional couplings to other degrees of freedom (see e.g. arXiv:1110.3327)

Model Independent Signatures

Observable 1: n_s and r

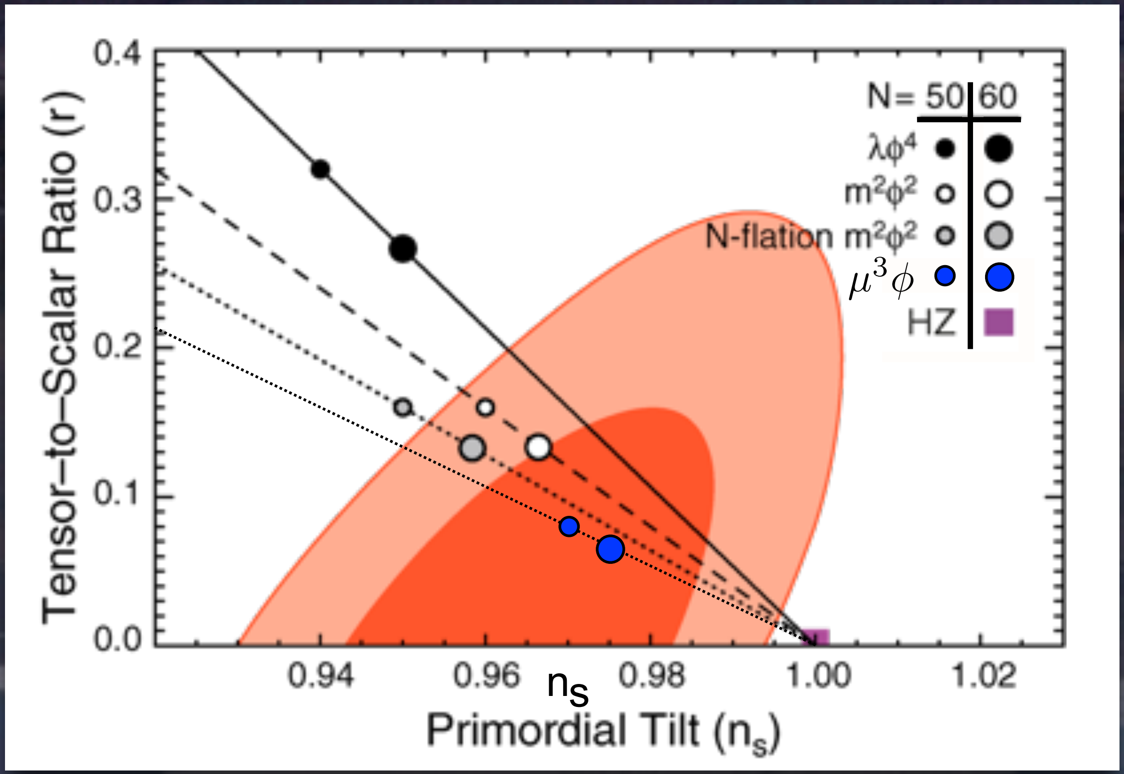
(Komatsu et al. 2010)



Model Independent Signatures

Observable 1: n_s and r

(modification of Komatsu et al. 2010)



Model Dependent Signatures I

Oscillations in the primordial power spectrum

In the presence of instanton corrections, the power spectrum gets modified.

This modification is not captured by the slow-roll approximation for the power spectrum because of parametric resonance, and the Mukhanov-Sasaki equation has to be solved.

Model Dependent Signatures I

Oscillations in the primordial power spectrum

$$\frac{d^2\mathcal{R}_k}{dx^2} - \frac{2(1+2\epsilon+\delta)}{x}\frac{d\mathcal{R}_k}{dx} + \mathcal{R}_k = 0$$

with

$$\epsilon = \epsilon_* - 3bf\sqrt{2\epsilon_*}\cos\left(\frac{\phi_k + \sqrt{2\epsilon_*}\ln x}{f}\right)$$

$$\delta = \delta_* - 3b\sin\left(\frac{\phi_k + \sqrt{2\epsilon_*}\ln x}{f}\right)$$

Model Dependent Signatures I

Oscillations in the primordial power spectrum

$$\frac{d^2\mathcal{R}_k}{dx^2} - \frac{2(1+\delta_{\rm osc}(x))}{x}\frac{d\mathcal{R}_k}{dx} + \mathcal{R}_k = 0$$

Look for a solution

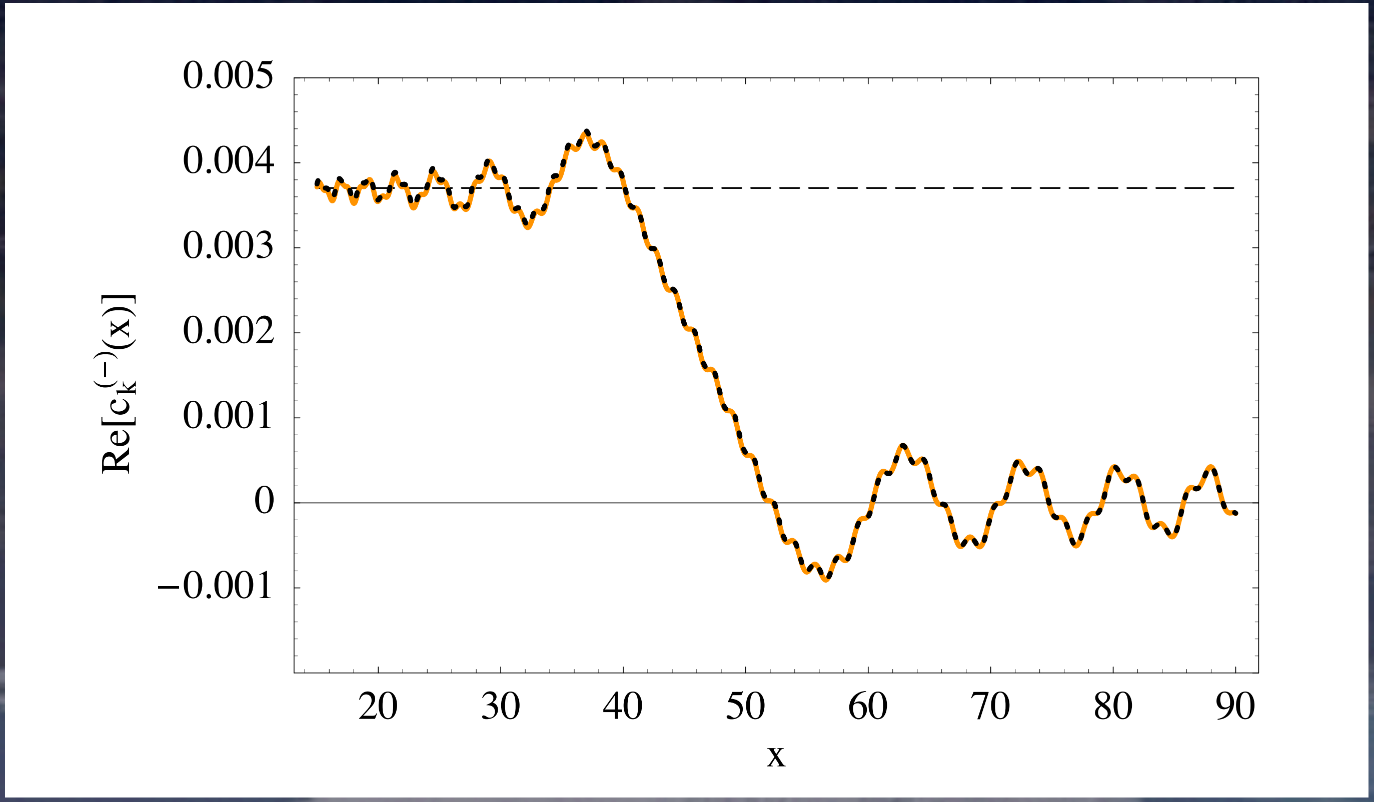
$$\mathcal{R}_k(x) = \mathcal{R}_{k,0}^{(o)}\left[i\sqrt{\frac{\pi}{2}}x^{3/2}H_{3/2}^{(1)}(x) - c_k^{(-)}(x)i\sqrt{\frac{\pi}{2}}x^{3/2}H_{3/2}^{(2)}(x)\right]$$

Then for large x

$$\frac{d}{dx}\left[e^{-2ix}\frac{d}{dx}c_k^{(-)}(x)\right] = -2i\frac{\delta_{\rm osc}(x)}{x}$$

Model Dependent Signatures I

Oscillations in the primordial power spectrum



Model Dependent Signatures I

One finds

$$\Delta^2_{\mathcal{R}}(k) = \Delta^2_{\mathcal{R}}(k_*) \left(\frac{k}{k_*}\right)^{n_s - 1} \left[1 + \delta n_s \cos\left(\frac{\phi_k}{f}\right)\right]$$

with

$$\delta n_s = \frac{12b}{\sqrt{1 + (3f\phi_*)^2}} \sqrt{\frac{\pi}{8} \coth\left(\frac{\pi}{2f\phi_*}\right) f\phi_*}$$

or for $f\phi_* \ll 1$:

$$\delta n_s = 3b(2\pi f\phi_*)^{1/2}$$

Model Dependent Signatures I

Constraints from WMAP5

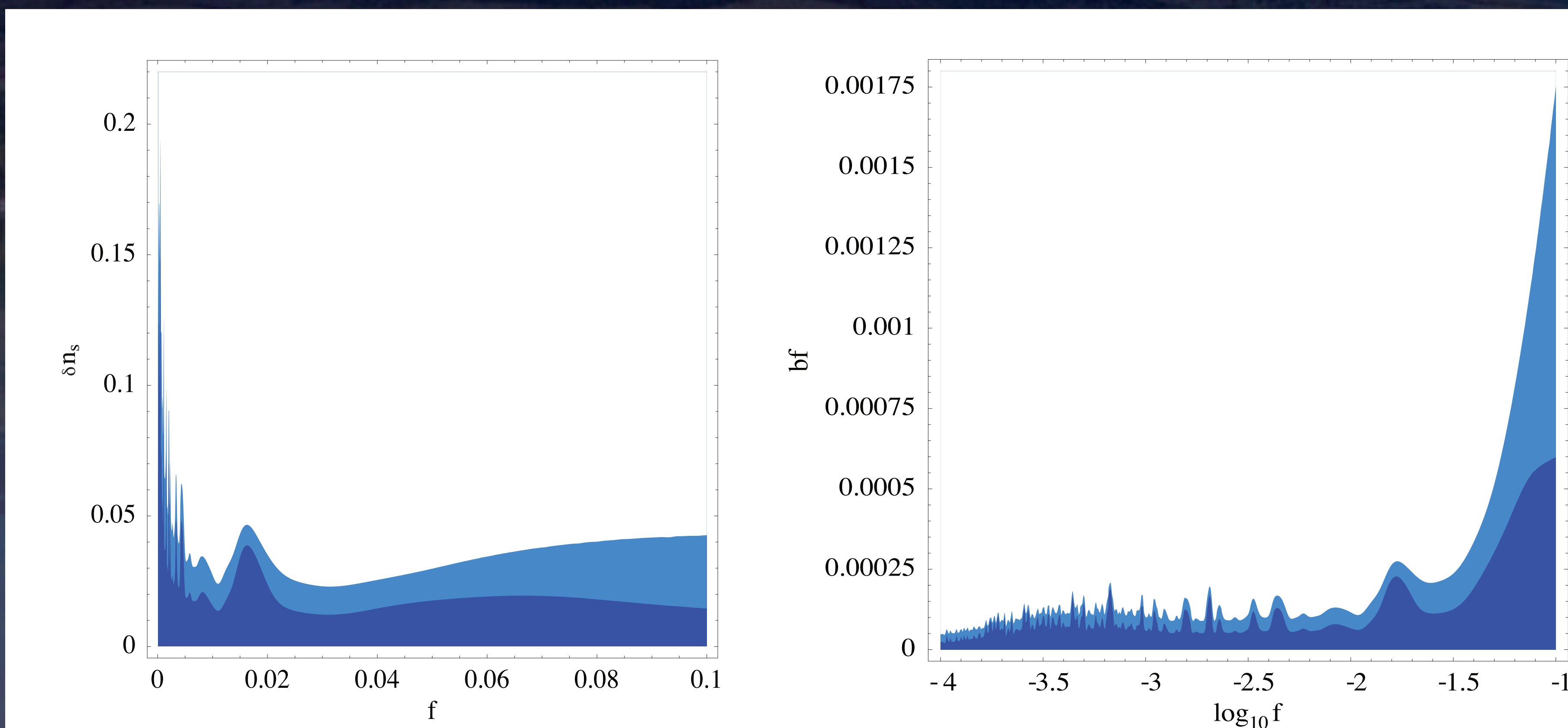
	Min	Max	Points
$\Omega_b h^2$	0.0212	0.0266	16
f	0.00009	0.1	512
δn_s	0	0.44	128
$\Delta\varphi$	$-\pi$	π	32

33 million spectra

(see Meerburg et al. arXiv:1109.5264 for WMAP7, and Huang et al. arXiv:1201.5955 for Planck and LSS forecasts)

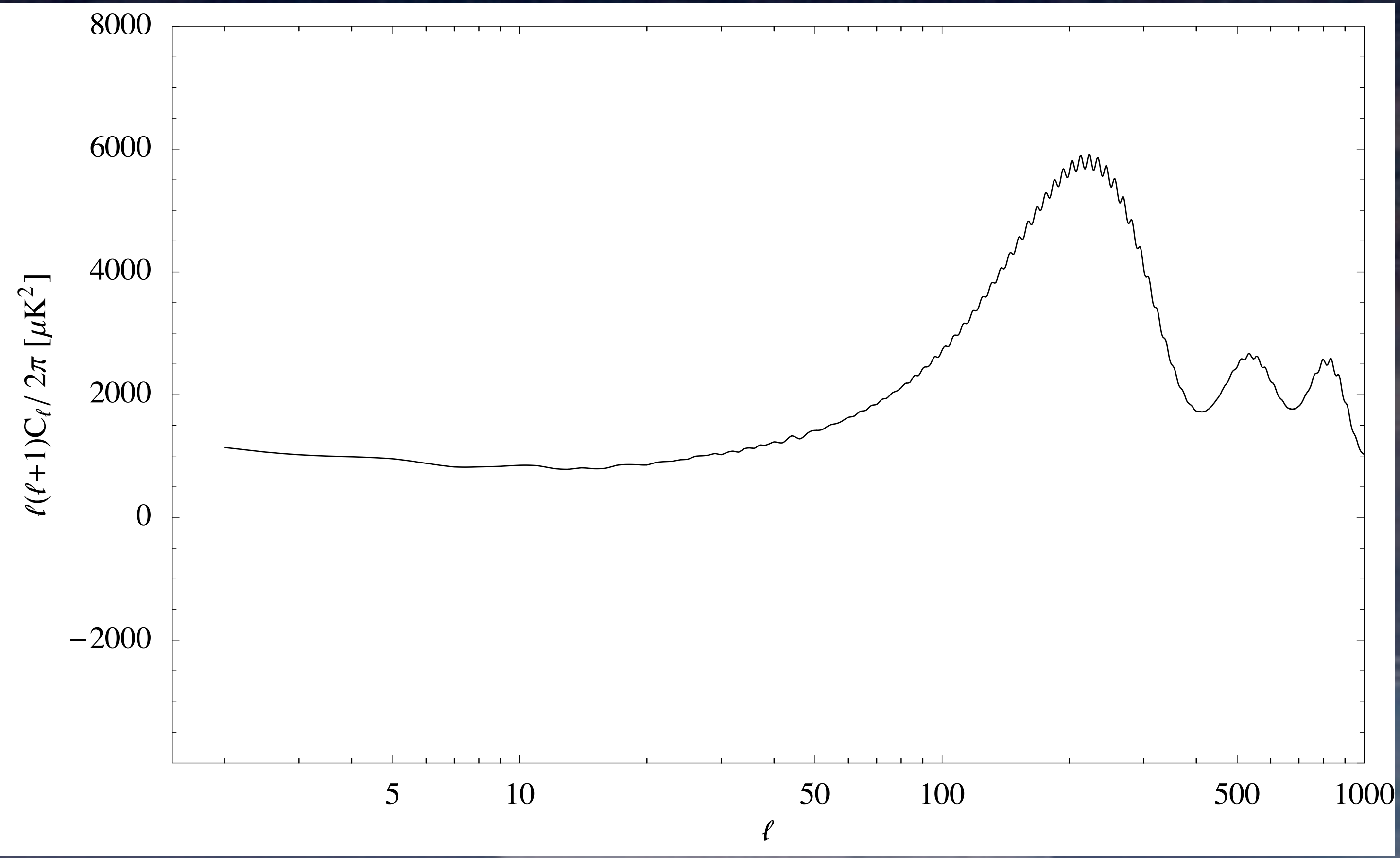
Model Dependent Signatures I

Constraints from WMAP5



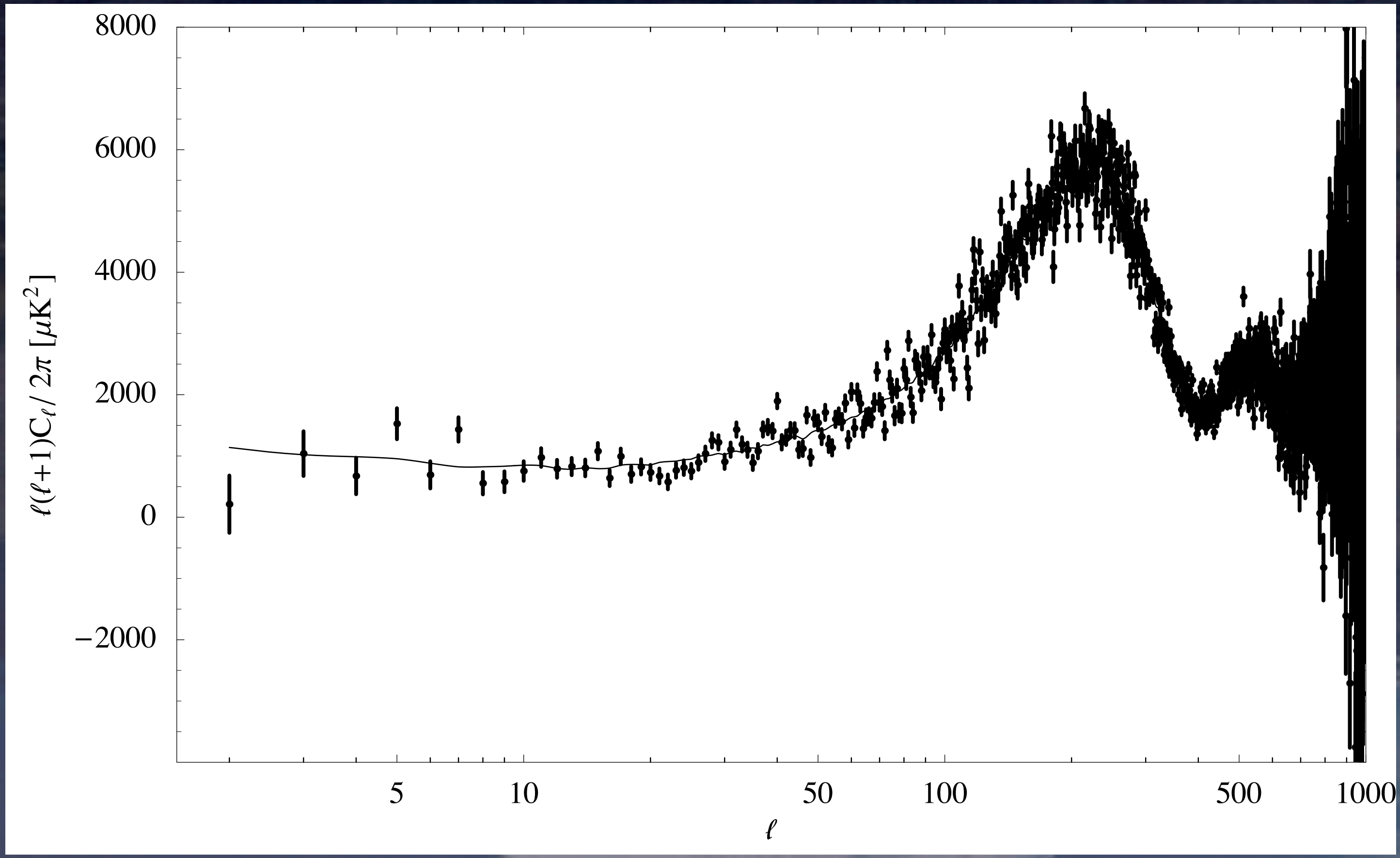
Model Dependent Signatures I

Constraints from WMAP5



Model Dependent Signatures I

Constraints from WMAP5



Model Dependent Signatures II

Resonant Non-Gaussianity

Models with large $\dot{\delta}$ can lead to large non-Gaussianities

(Chen, Easther, Lim 2008)

$$\langle \mathcal{R}(\mathbf{k_1},t)\mathcal{R}(\mathbf{k_2},t)\mathcal{R}(\mathbf{k_3},t)\rangle =$$

$$-i\int_{-\infty}^{t} dt' \langle [\mathcal{R}(\mathbf{k_1},t)\mathcal{R}(\mathbf{k_2},t)\mathcal{R}(\mathbf{k_3},t), H_I(t')]\rangle$$

with

$$H_I(t) \supset -\int d^3x\, a^3(t)\epsilon(t)\dot{\delta}(t)\mathcal{R}^2(\mathbf{x},t)\dot{\mathcal{R}}(\mathbf{x},t)$$

Model Dependent Signatures II

Observable III: Resonant Non-Gaussianity

After some algebra

$$\frac{\mathcal{G}(k_1,k_2,k_3)}{k_1k_2k_3} = \frac{1}{8}\int_0^\infty dX\,\frac{\dot{\delta}}{H}e^{-iX}$$

$$\left[-i - \frac{1}{X}\sum_{i\neq j}\frac{k_i}{k_j} + \frac{i}{X^2}\frac{K(k_1^2+k_2^2+k_3^2)}{k_1k_2k_3}\right] + c.c$$

Model Dependent Signatures II

Observable III: Resonant Non-Gaussianity

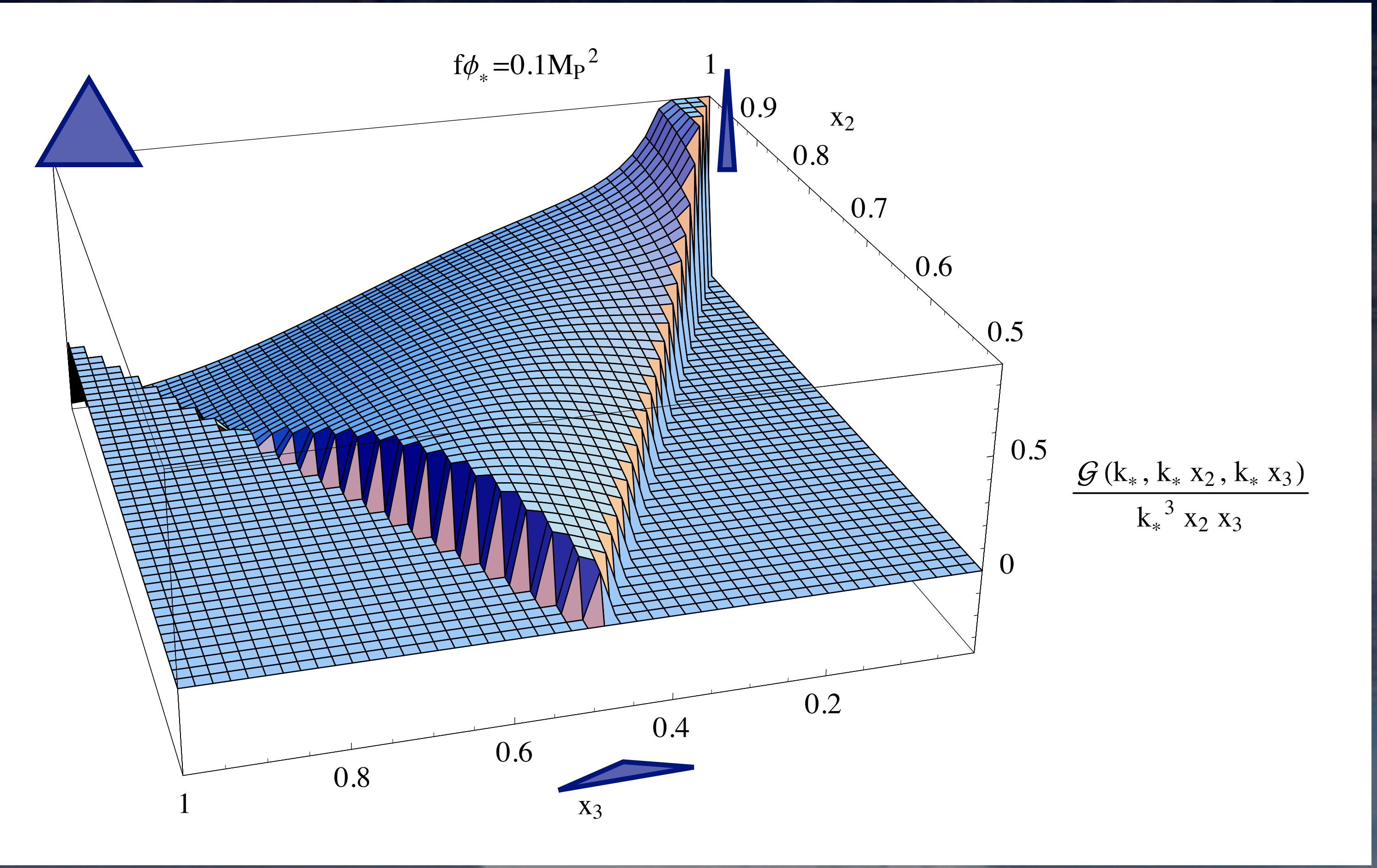
$$\frac{\mathcal{G}(k_1,k_2,k_3)}{k_1k_2k_3} = f^{\rm res}\left[\sin\left(\frac{\ln K/k_*}{f\phi_*}\right) + f\phi_*\sum_{i\neq j}\frac{k_i}{k_j}\cos\left(\frac{\ln K/k_*}{f\phi_*}\right)\right]$$

with $$K = k_1 + k_2 + k_3$$

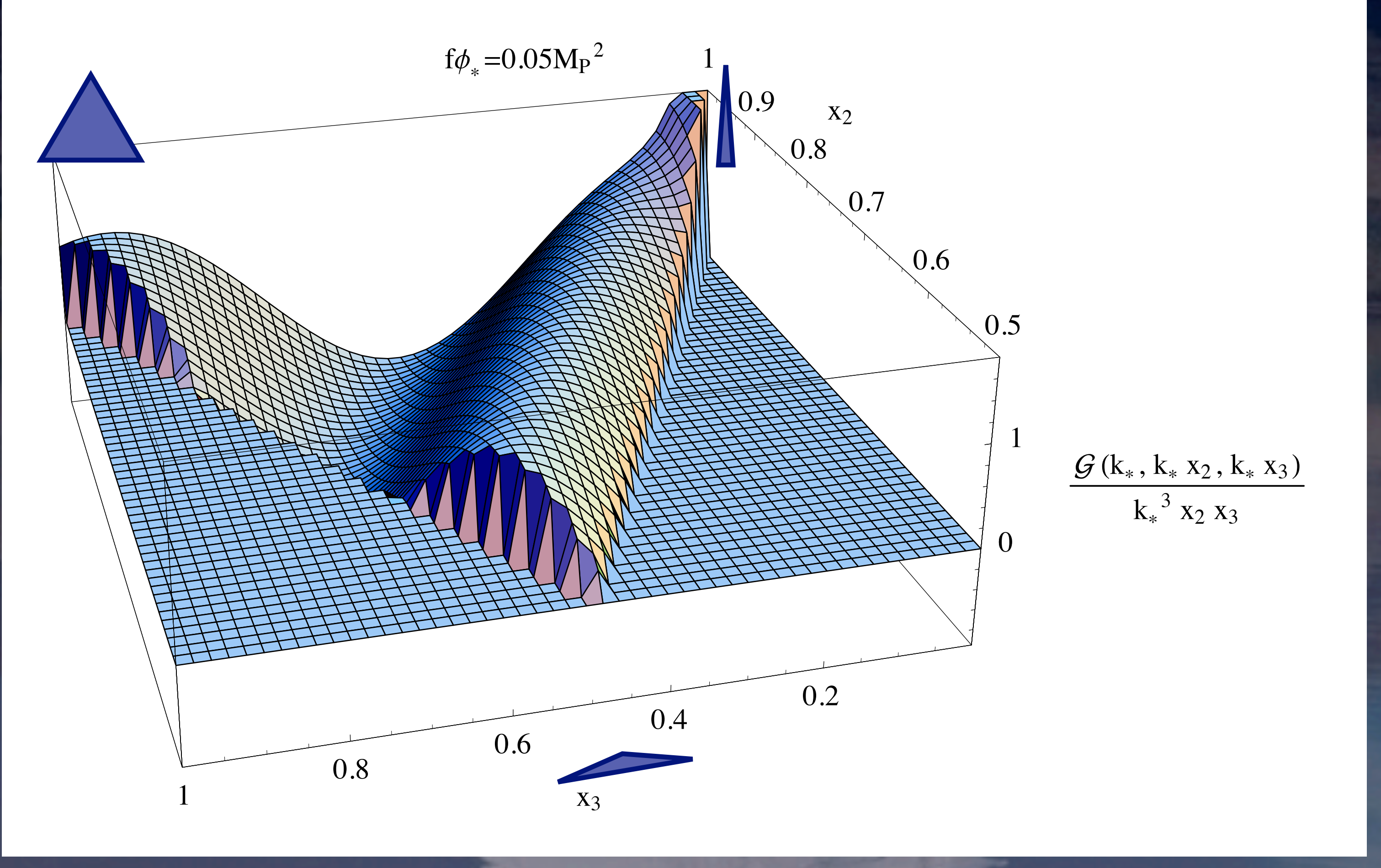
$$f^{\rm res} = \frac{3\sqrt{2\pi}b}{8(f\phi_*)^{3/2}}$$

(This satisfies the consistency condition.)

Model Dependent Signatures II



Model Dependent Signatures II



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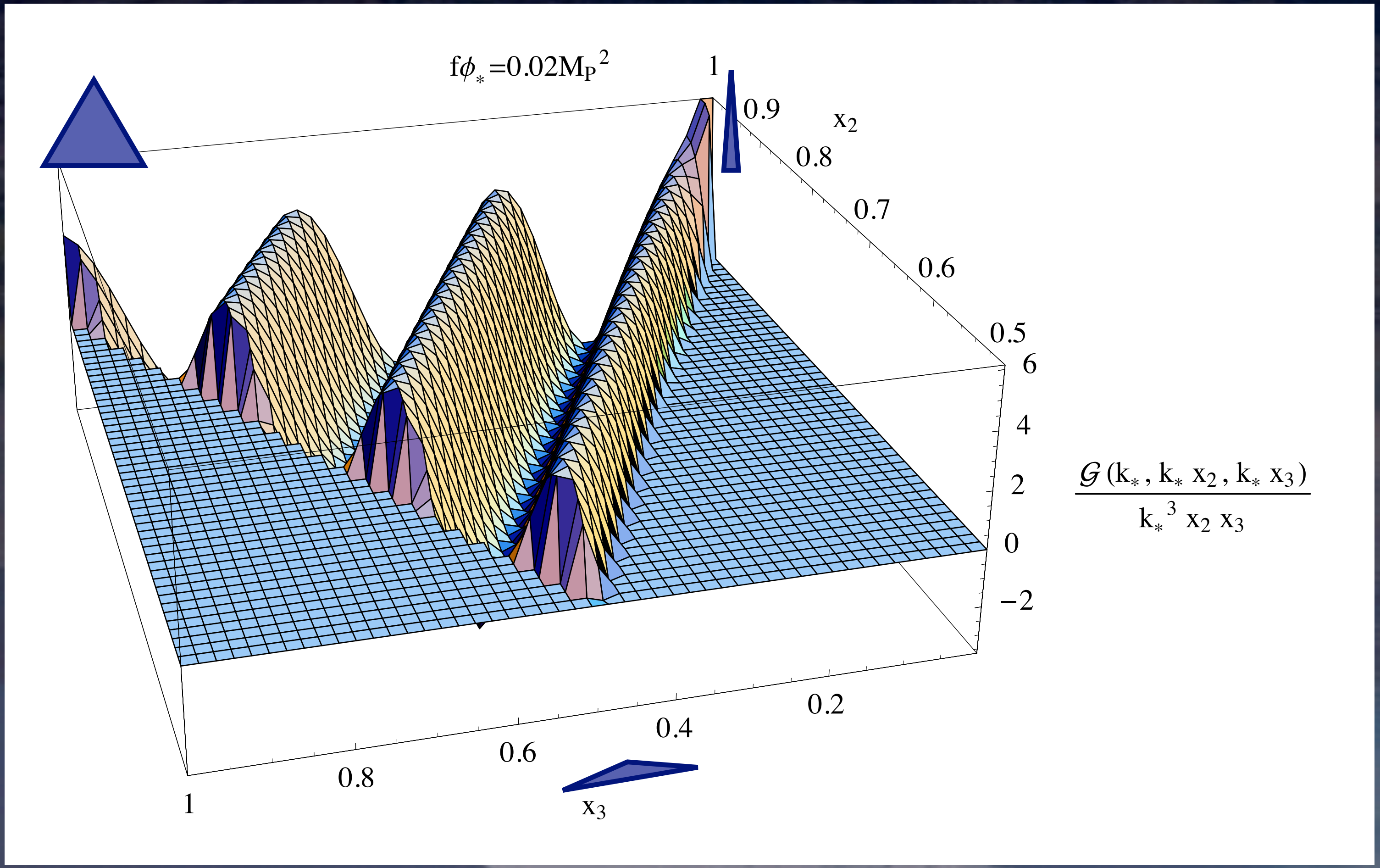
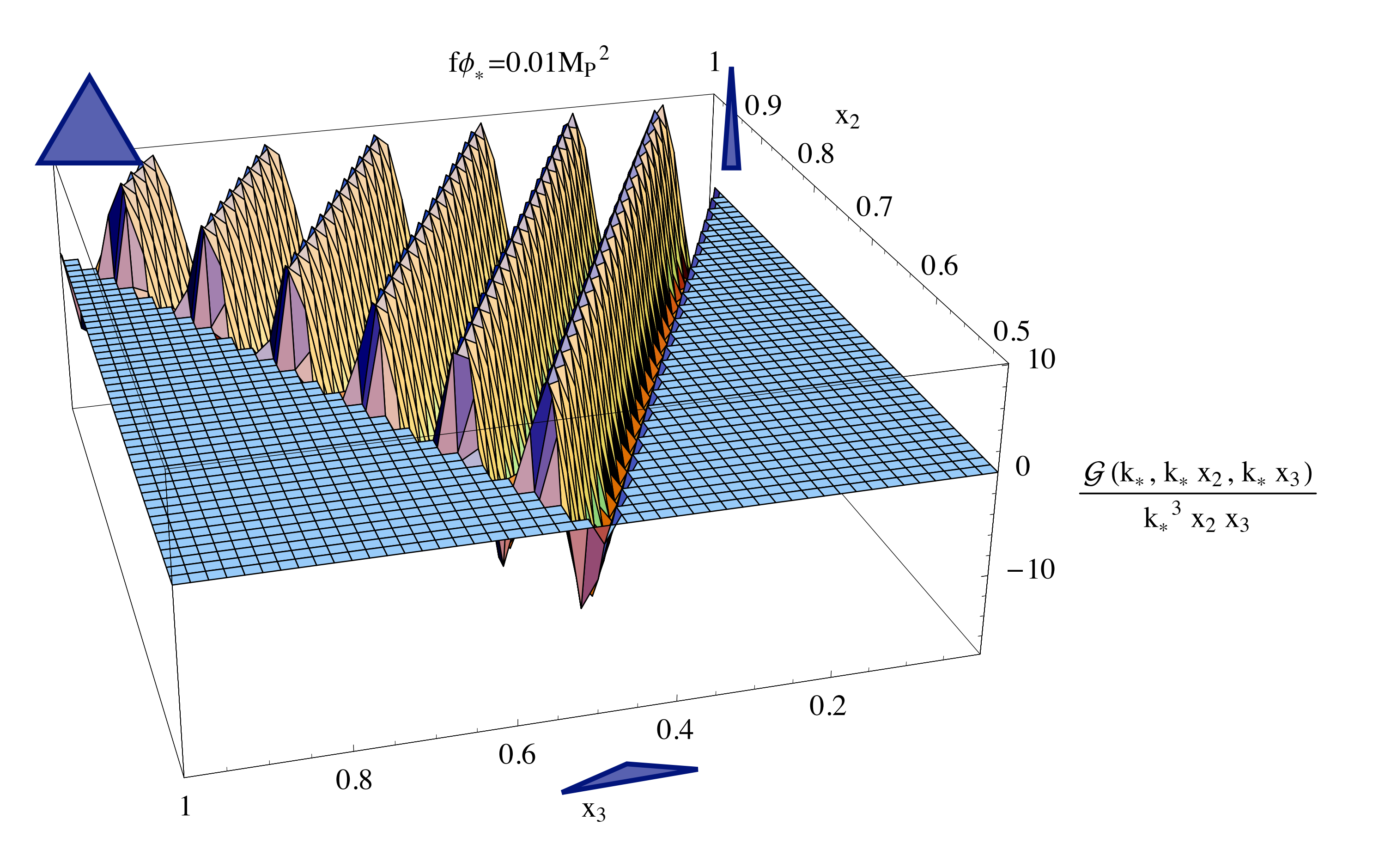
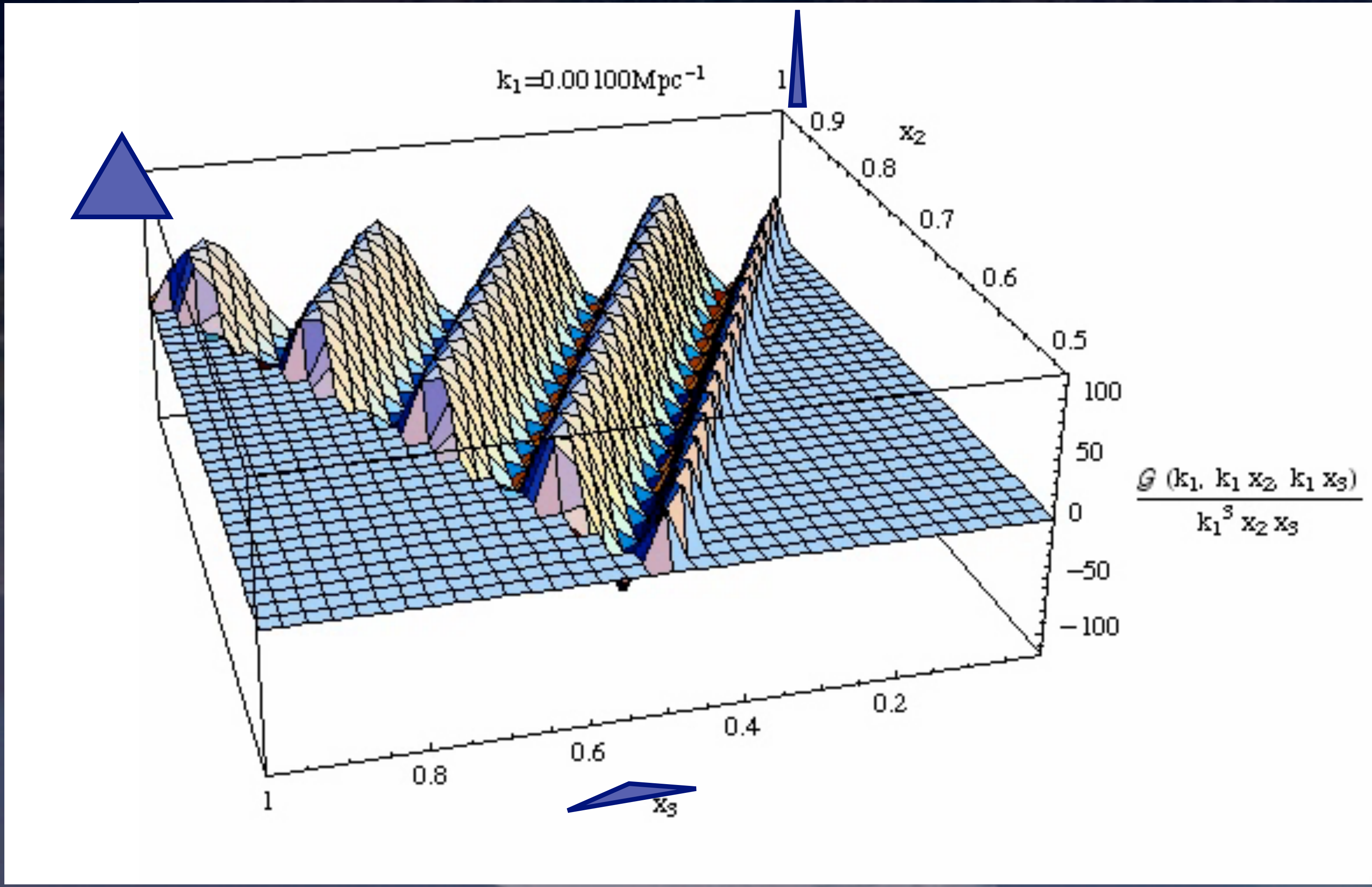


Model Dependent Signatures II



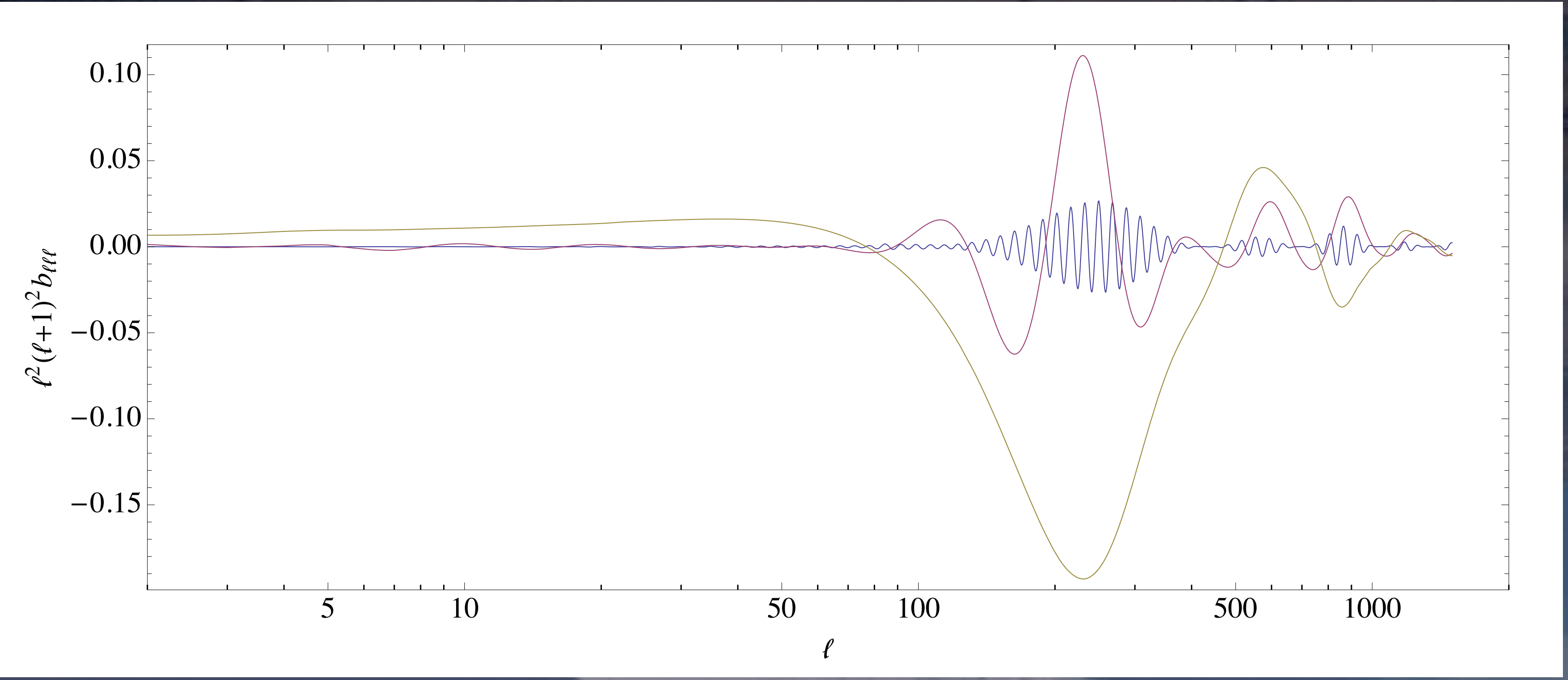
Model Dependent Signatures II



Model Dependent Signatures II

If one takes the best-fit point of the 2-pt analysis seriously, one expects a signal with f_{res}~400.

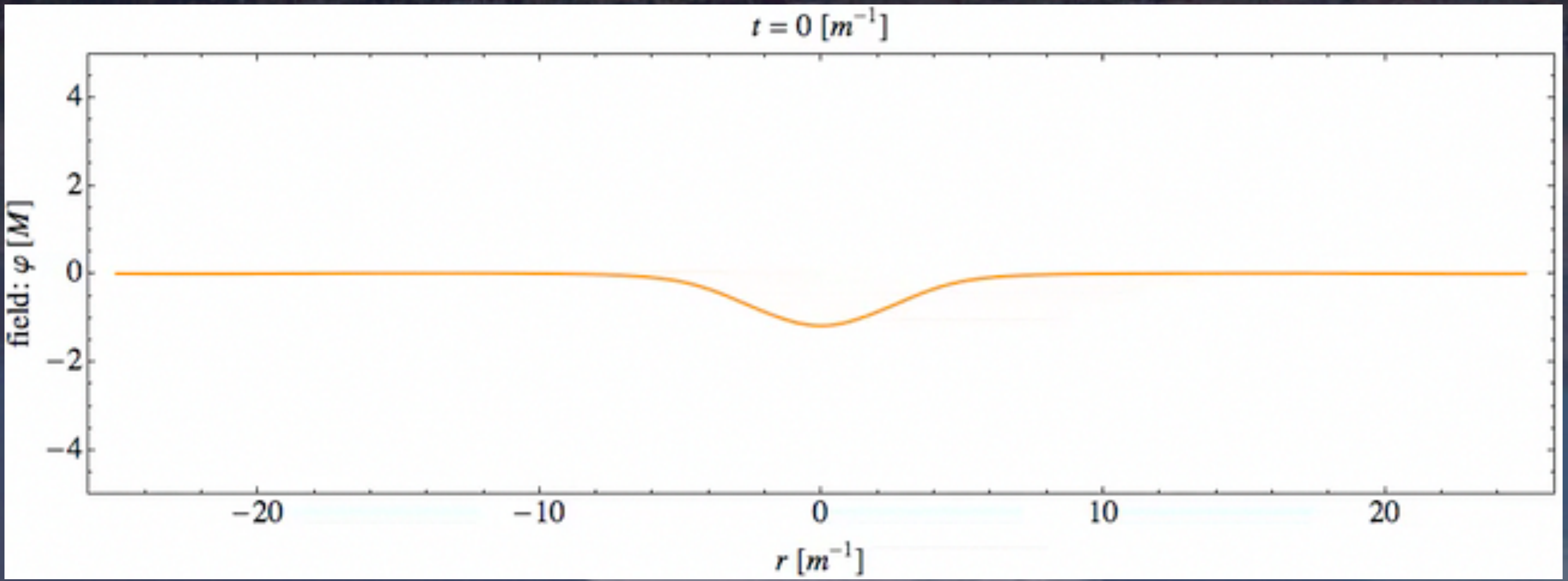
Whether this is detectable remains to be seen.



Oscillons

What are they?

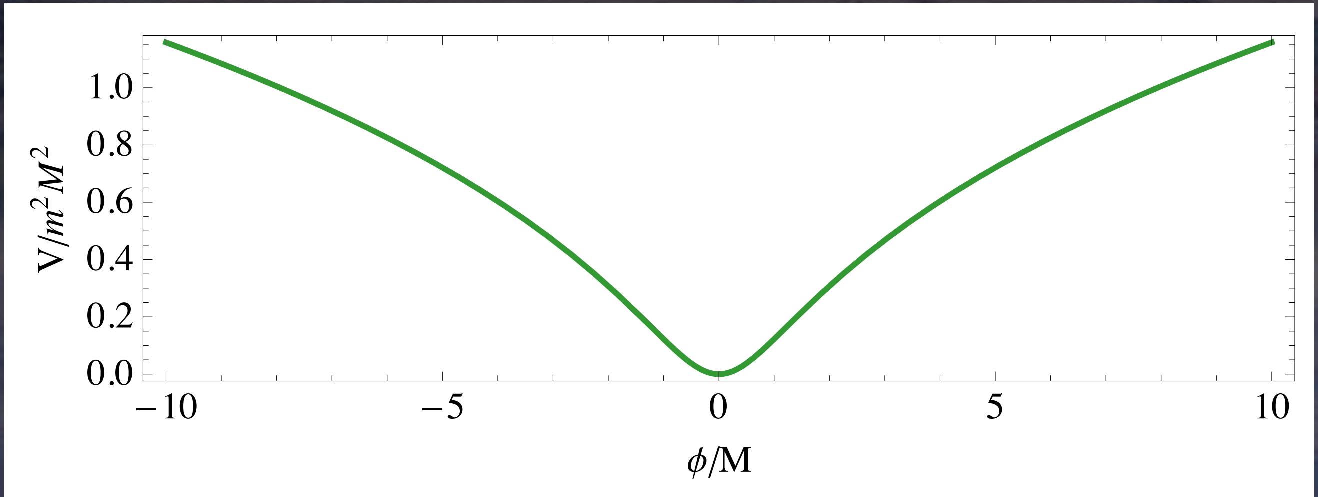
Oscillons are long-lived, localized, oscillatory configurations of a scalar field



Oscillons

Under what circumstances do they exist?

They can only exist if the potential flattens out away from the minimum.



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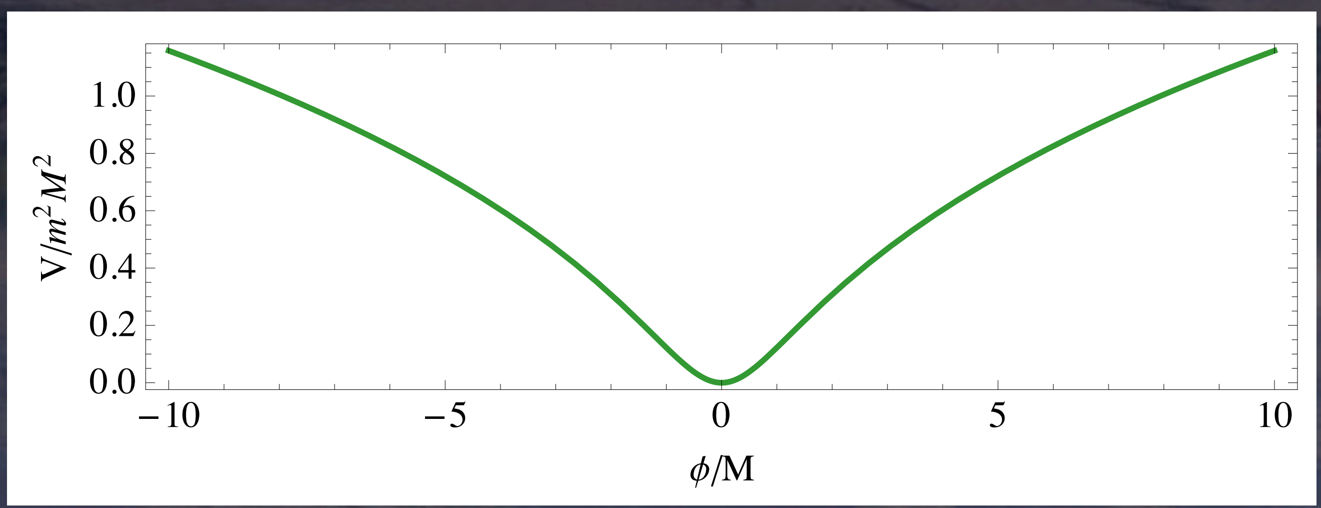

$$\underbrace{(-\omega^2 + m^2)\phi}_{>0} \underbrace{- \nabla^2\phi}_{>0} + V' - m^2\phi = 0$$

Oscillons

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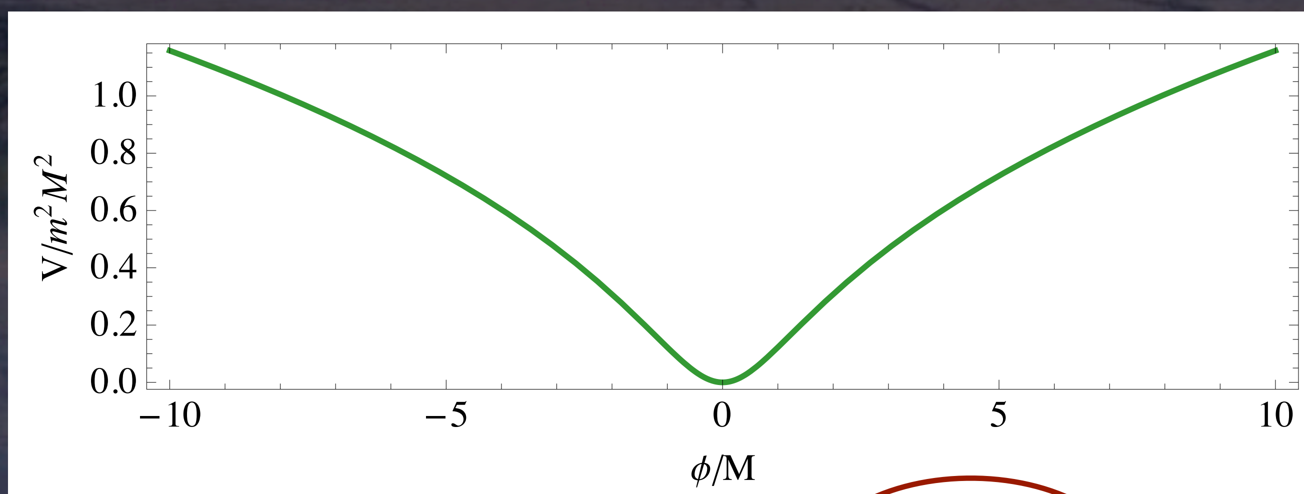


$$\underbrace{(-\omega^2 + m^2)\phi}_{>0} \underbrace{- \nabla^2\phi}_{>0} + \underbrace{V' - m^2\phi}_{<0} = 0$$

Oscillons

Longevity

Oscillons classically radiate at an exponentially small rate.

Segur, Kruskal, Phys. Rev. Lett. 58 (1987)

While there is no conserved charge (Q-balls) associated with them, there is an adiabatic invariant.

Kasuya, Kawasaki, Takahashi, hep-ph/0209358

Quantum effects turn the rate into a power law.

Hertzberg, arXiv:1003.3459

Oscillons

Stability

Small oscillons suffer an instability in the presence of perturbations comparable to their size.

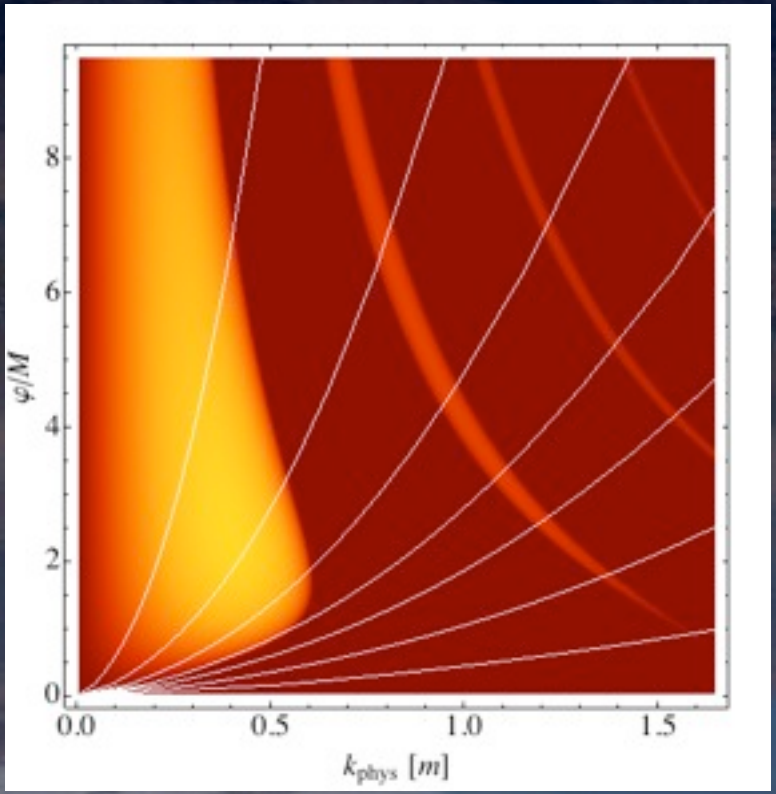
Very wide oscillons suffer an instability in the presence perturbations much smaller than their size.

e.g. Amin, Shirokoff, arXiv:1002.3880

Oscillons

Formation at the end of inflation

Parametric resonance in the equations for the perturbations leads to an instability which causes the inflaton to fragment and form oscillons.

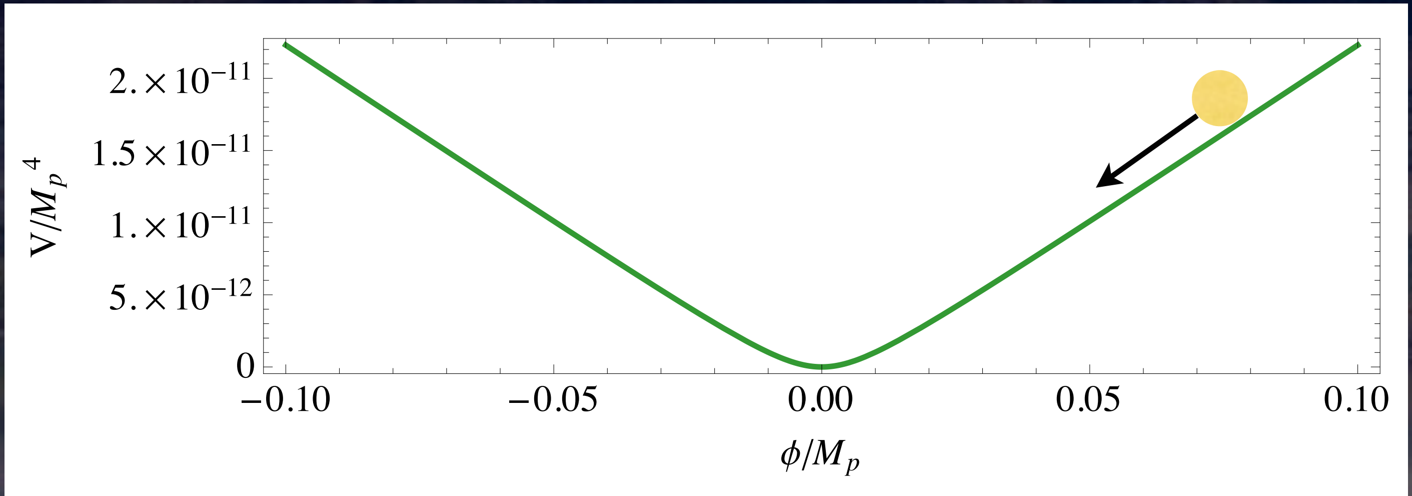


$$\delta\ddot{\phi} + 3H\delta\dot{\phi} - \frac{1}{a^2}\nabla^2\delta\phi + V''(\bar{\phi})\delta\phi$$

$$+16\pi G\frac{\dot{\bar{\phi}}}{H}V'(\bar{\phi})\delta\phi + (8\pi G)^2\frac{\dot{\bar{\phi}}^2}{H^2}V(\bar{\phi})\delta\phi = 0$$

Oscillons

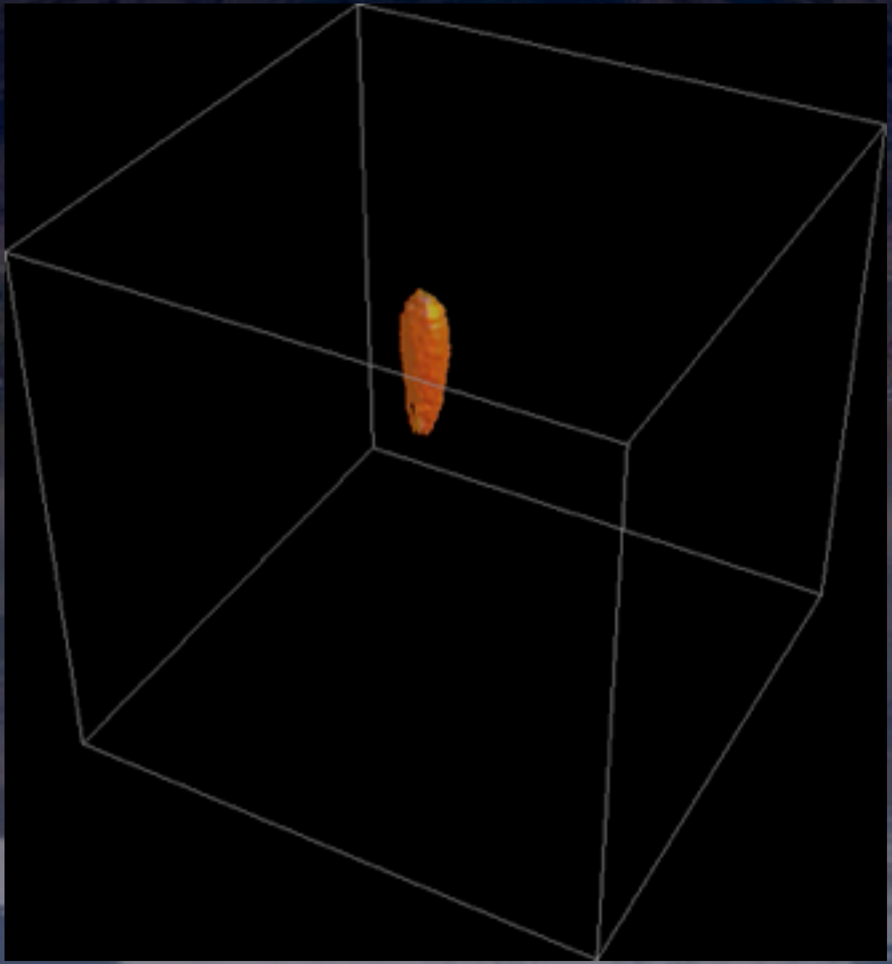
Formation at the end of inflation



The numerical calculation begins at the first turning point. It assume that the couplings to the degrees of freedom on the brane are small enough so that $\Gamma \ll H$ at the end of inflation.

Oscillons

Oscillon formation in the monodromy potential



PSpectRe

Conclusions

- String theory seems to contain the right ingredients to realize large field inflation.
- The scenario has interesting signatures:
  A large tensor to scalar ratio, potentially a modulated temperature anisotropy spectrum as well as resonant non-Gaussianities.
- Resonant non-Gaussianity is currently poorly constrained and deserves further study independent of the stringy scenario.

Conclusions

- In these models oscillons may form at the end of inflation. The properties of oscillons, too, deserve further study both in the context of the stringy model and in their own right.
- More explicit geometries are desirable, reheating in these models should be studied, ...


Thank you