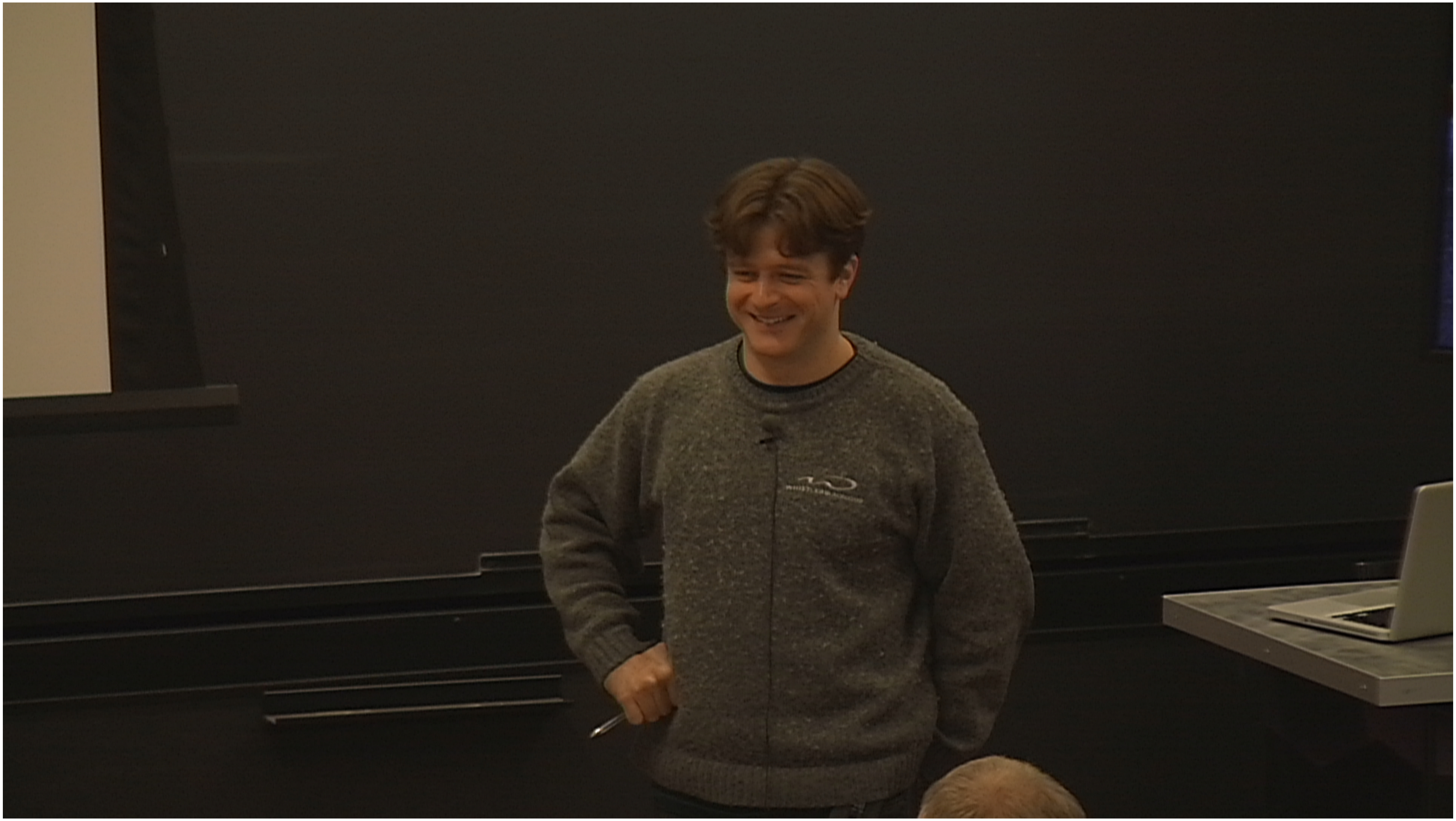
Title: Axion Monodromy and its Signatures

Date: Nov 22, 2011 11:00 AM

URL: http://pirsa.org/11110071

Abstract: The study of the anisotropies in the cosmic microwave background radiation over the past two decades has provided us with important information about the early universe. In particular, there is strong evidence that these anisotropies were generated long before the cosmic microwave radiation was emitted. The most commonly studied idea is that they originated as quantum fluctuations during a period of inflation. In addition to a spectrum of scalar perturbations consistent with the one that has been observed, inflation also predicts the presence of gravitational waves. These might be observable in the polarization of the cosmic microwave background. An observation of this signal would indicate that the inflaton must have traversed a super-Planckian distance. Realizing this in string theory has been challenging. I will describe the basic ingredients for a string theoretic setup in which the inflaton can move over a super-Planckian distance, leading to an observable gravitational wave signal within string theory. In addition to an observable tensor signal, the model may also lead to other interesting signatures which I will discuss such as modulations in the power spectrum of scalar perturbations, interesting shapes of non-Gaussianities, and possibly the formation of oscillons at the end of inflation.



Axion Monodromy Inflation and its Signatures

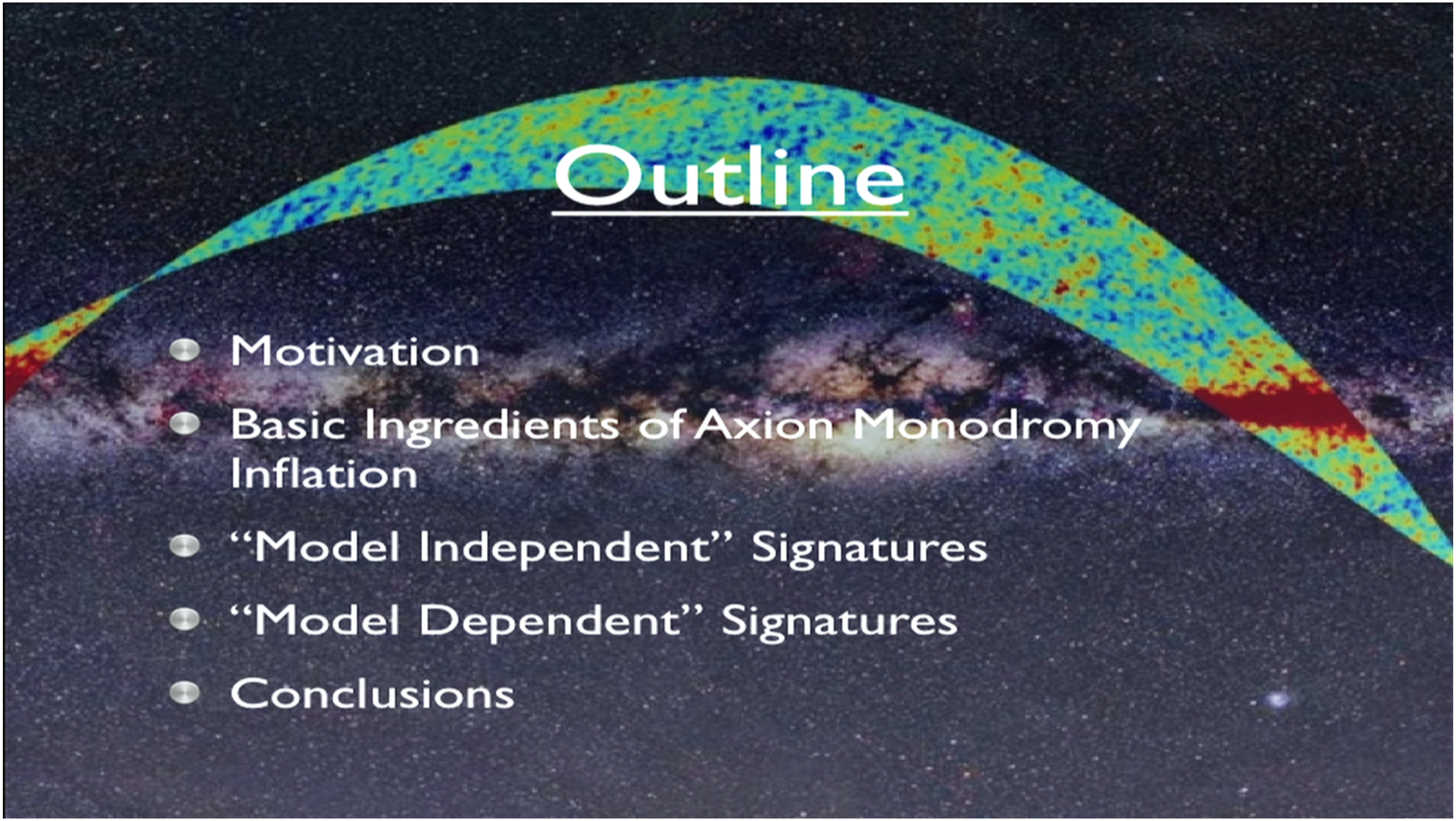
Raphael Flauger

arXiv:0907.2916
w/
Liam McAllister, Enrico Pajer, Alexander Westphal, Gang Xu

arXiv:1002.0833
w/
Enrico Pajer

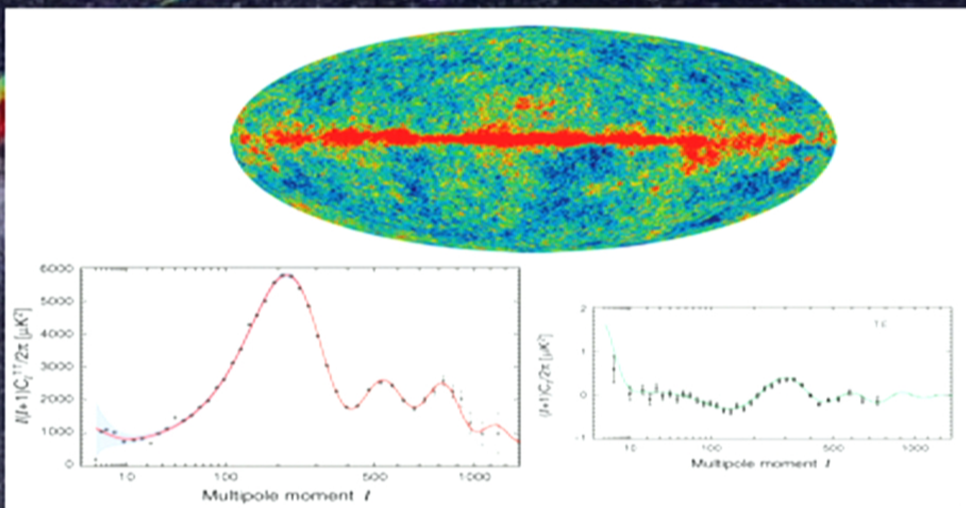
arXiv:1106.3335
w/
Mustafa Amin, Richard Easther, Hal Finkel, Mark Hertzberg

Cosmology and Gravitation, Perimeter Institute, November 22, 2011


Outline
Motivation
Basic Ingredients of Axion Monodromy Inflation
"Model Independent" Signatures
"Model Dependent" Signatures
Conclusions

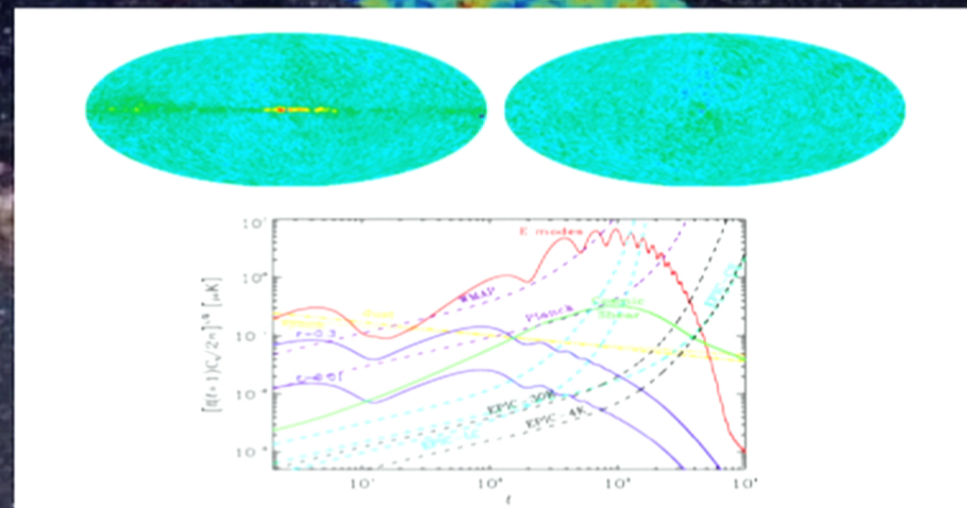
Motivation

(Jarosik et al. 2010)



(Larson et al. 2010)

(Jarosik et al. 2010)



(Bock et al. 2009)

Motivation

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)**

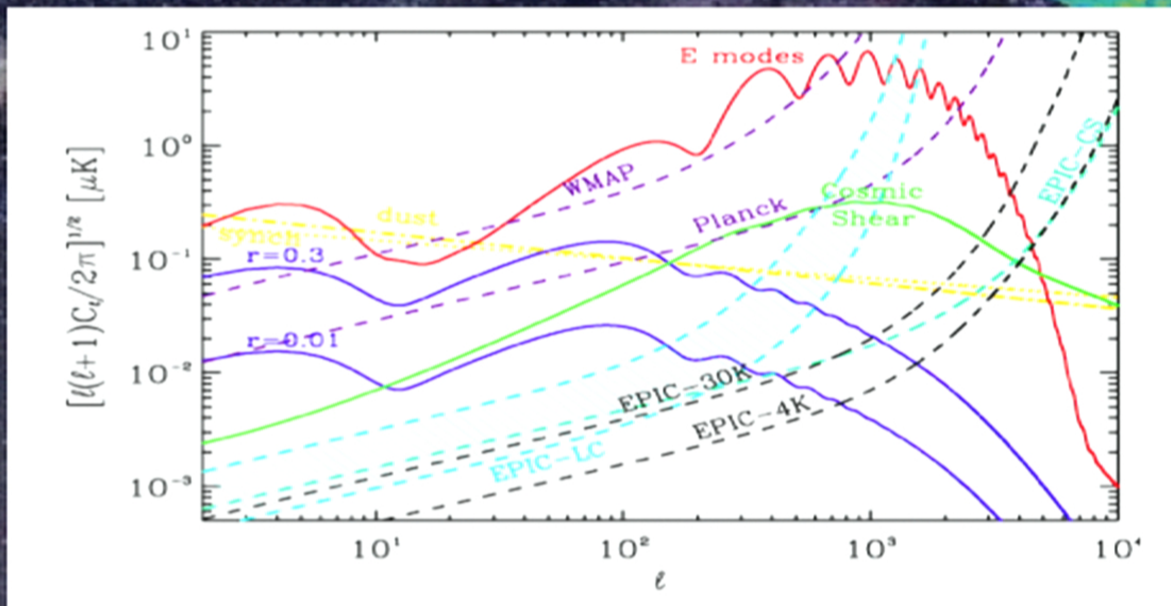
* For single field models with canonical kinetic term

Motivation

WMAP+BAO+H_0: $r < 0.24$ $(95\% \, C.L.)$ **(Komatsu et al. 2010)**

Future experiments: $r \sim 0.001$

(Bock et al. 2009)

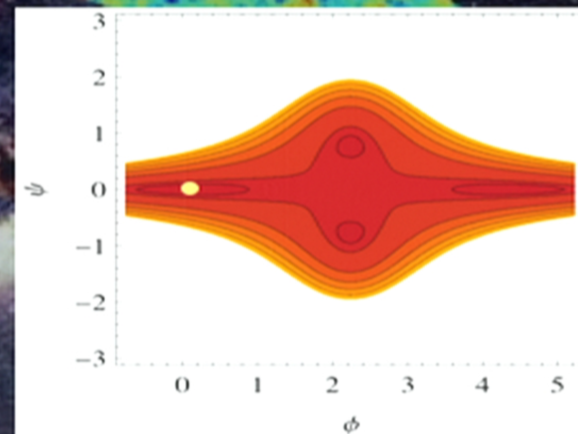


Motivation

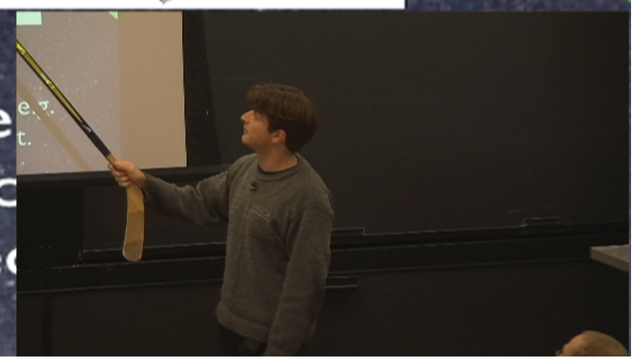
Motion of the scalar field over super-Planckian distance is hard to control in an effective field theory

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

$$(\Lambda < M_p)$$



The c_n are typically unknown.
Even if they were known, the effective
generically expected to break down fo
because other degrees of freedom bec



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 more recently

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

Berg, Pajer, Sjors, arXiv:0912.1341

interesting studies of monodromy in EFT

Kaloper, Lawrence, Sorbo, arXiv:1101.0026

Dubovsky, Lawrence, Roberts, arXiv:1105.3740


Basic Ingredients for Axion
Monodromy Inflation

Axion Monodromy Inflation

Consider string theory on $M \times X$

Axions arise from integrating gauge potentials over non-trivial cycles in the compactification manifold.

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

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

where $\Sigma_\alpha^{(p)}$ is an element of an integral basis of $H_p(X, \mathbb{Z})$

Axion Monodromy Inflation

These fields possess a shift symmetry to all orders in string perturbation theory.

The vertex operator for $b_I(x)$ in the limit of vanishing momentum is

$$V_{b_I}(0) = \int_{\mathcal{W}} d^2\xi \epsilon^{\alpha\beta} \partial_\alpha Y^i \partial_\beta Y^j \omega^I_{ij}(Y(\xi)) = \int_{\varphi(\mathcal{W})} \omega^I$$

with $\omega^I \in H^2(X, \mathbb{Z})$ dual to Σ_I

vanishes if $\varphi(\mathcal{W}) = \partial\mathcal{C}$ so that coupling vanishes.

Axion Monodromy Inflation

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

Breaking by branes

This implies the following potential

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

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

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

Quadratic for small field values, linear for large field values

Axion Monodromy Inflation

Inflationary Potential

For large field values in terms of the canonically normalized fields the potential then becomes

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

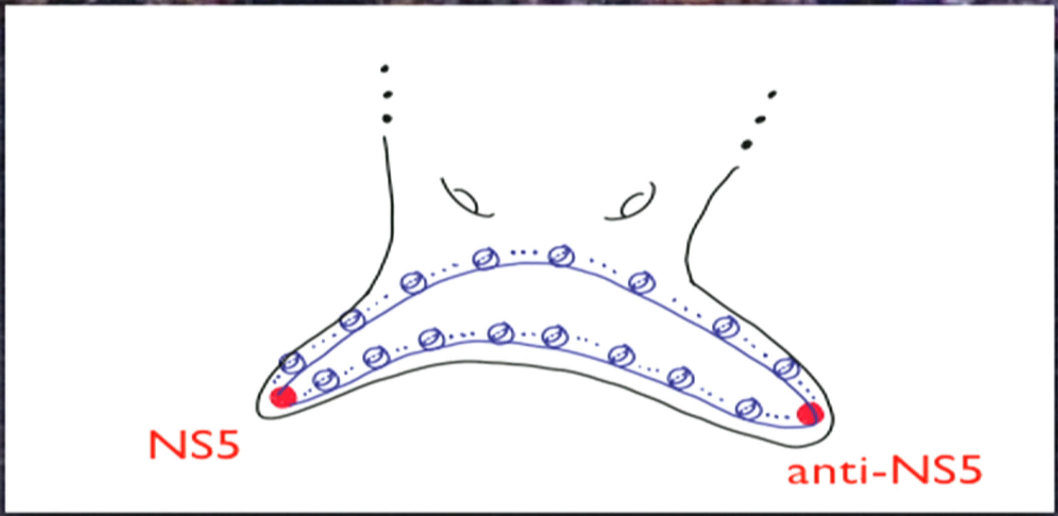
with $$\mu = \frac{\epsilon^{1/3}(2\pi)^3 g_s}{L^{10/3}} M_p \qquad \text{for b}$$

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

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.

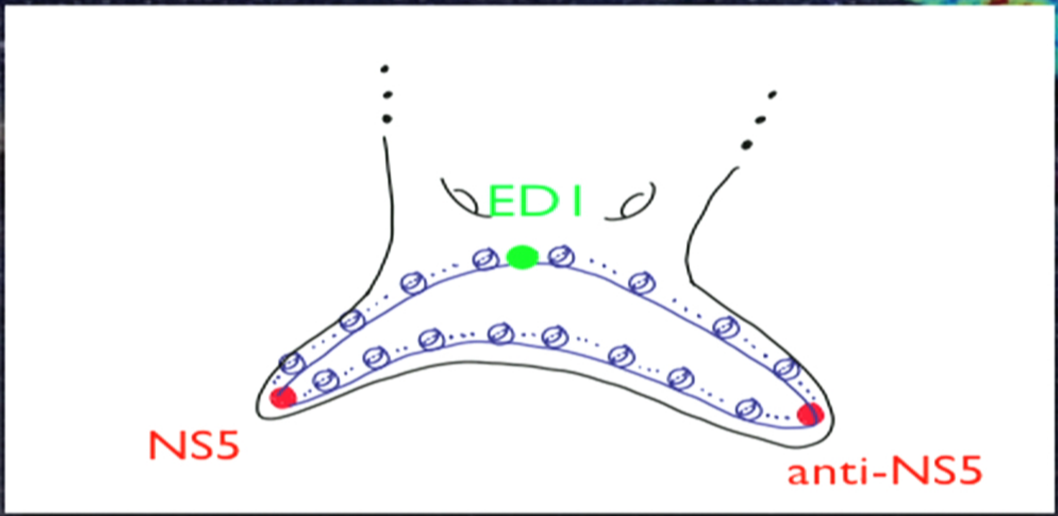
The backreaction on the geometry must be controlled.

Higher derivative corrections must be negligible.

Instanton corrections must be controlled.

Axion Monodromy Inflation

Instanton corrections may lead to interesting signatures.



$$K = -2\log(\mathcal{V}_E + e^{-S_{ED1}}\cos(c))$$

Signatures

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

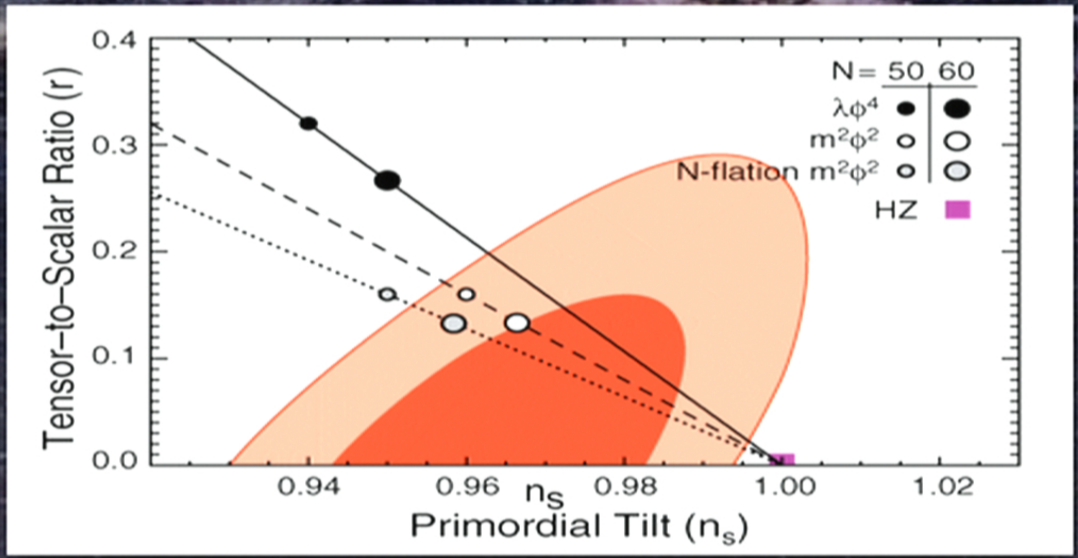
$$V(\phi) = \mu^3 \phi + b\mu^3 f \cos(\phi/f)$$

possibly with additional couplings to other degrees of freedom

Model Independent Signatures

Observable 1: n_s and r

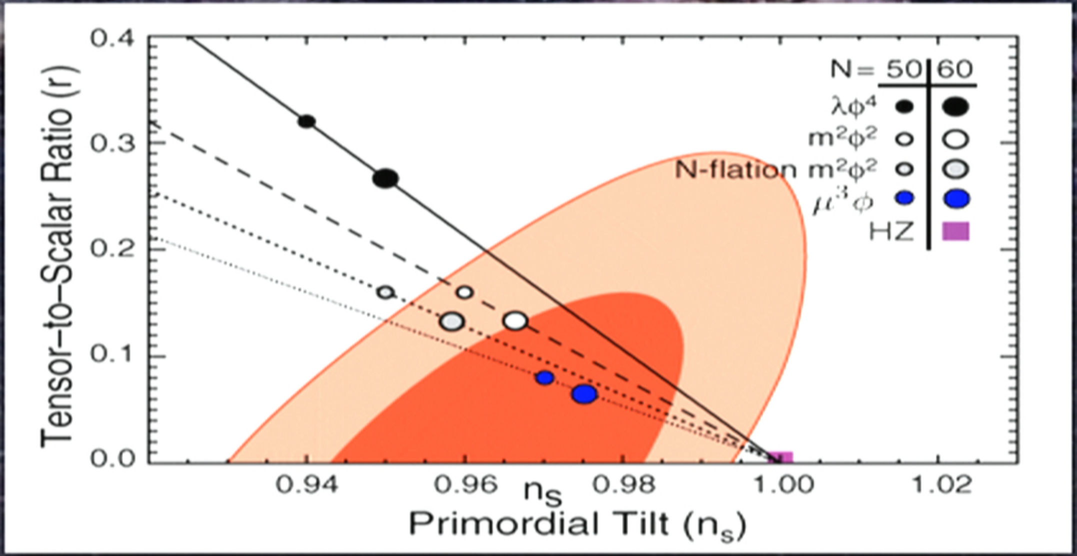
(Komatsu et al. 2010)



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

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Model Dependent Signatures I

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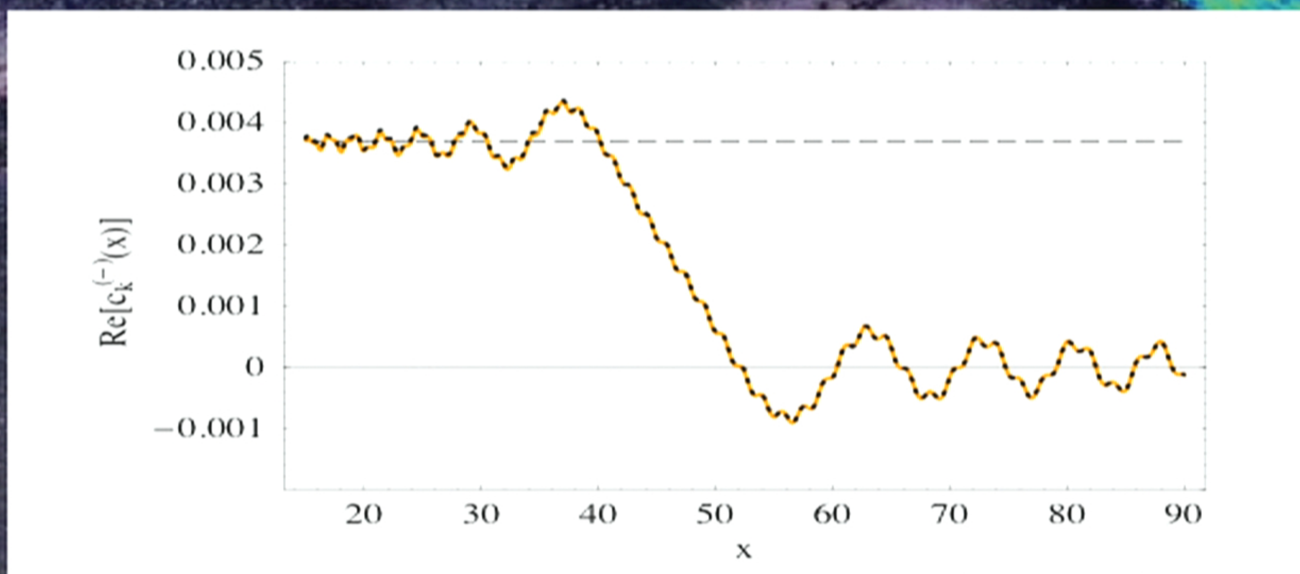


Model Dependent Signatures I

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

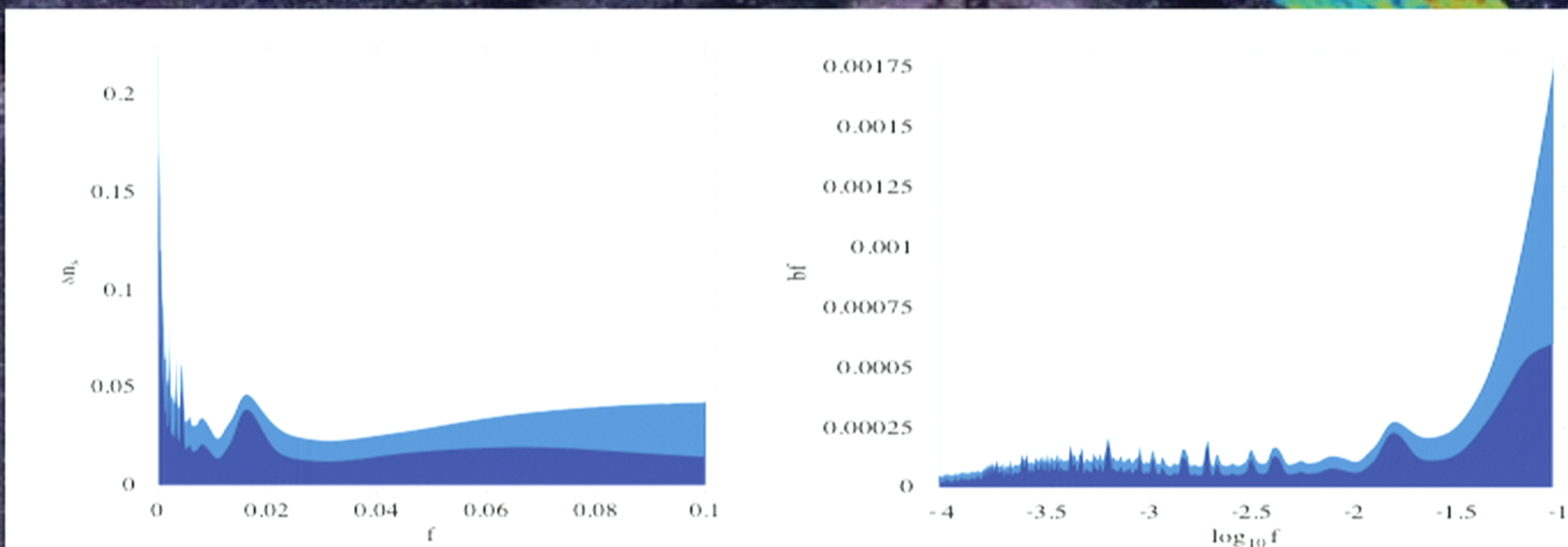
Constraints from WMAP5

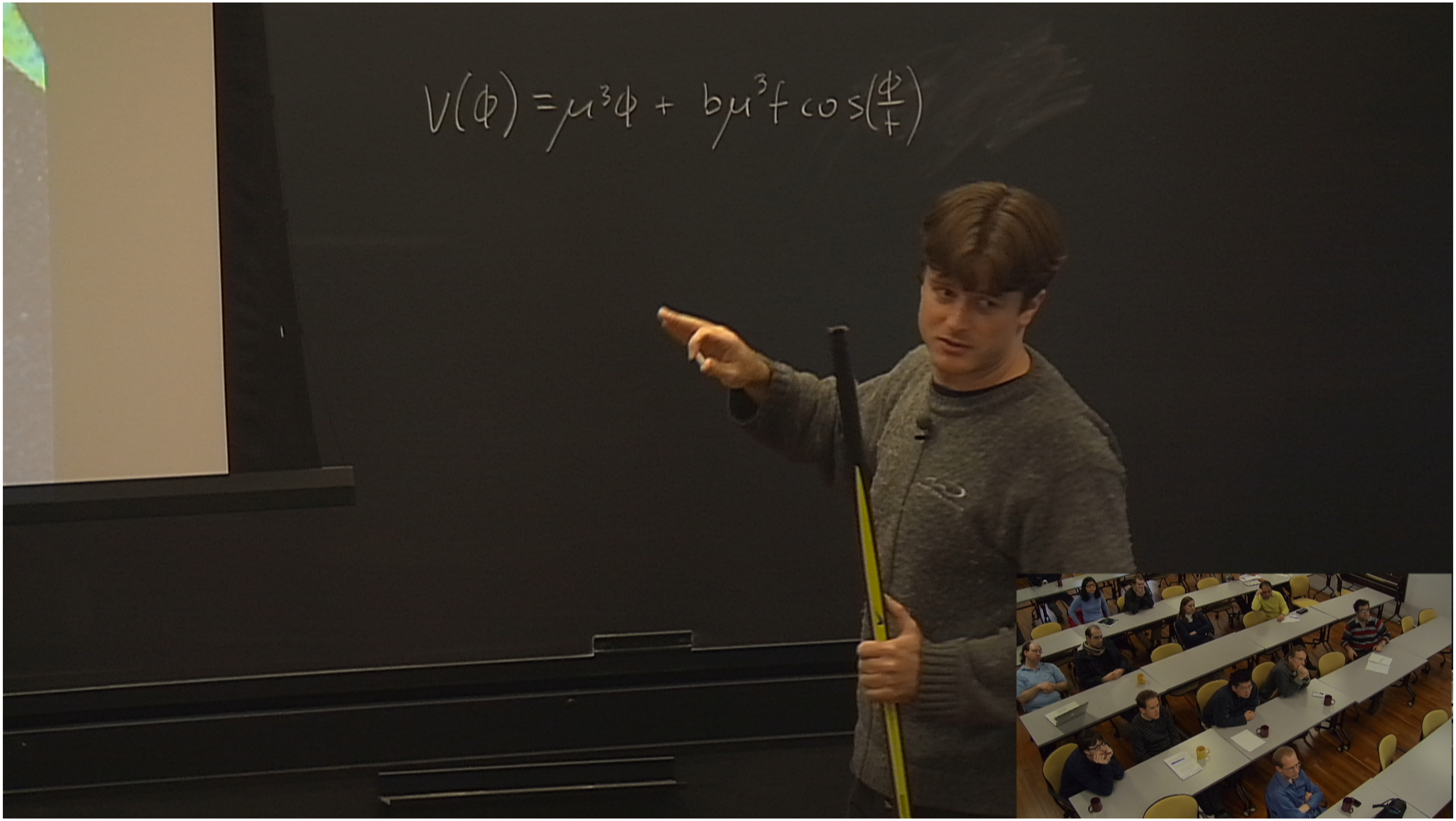
	Min	Max	Points
$\Omega_b h^2$	0.0212	0.0266	16
f	0.00009	0.1	512
δn_s	0	0.44	128
$\Delta\varphi$	$-\pi$	π	32

33 million spectra

Model Dependent Signatures I

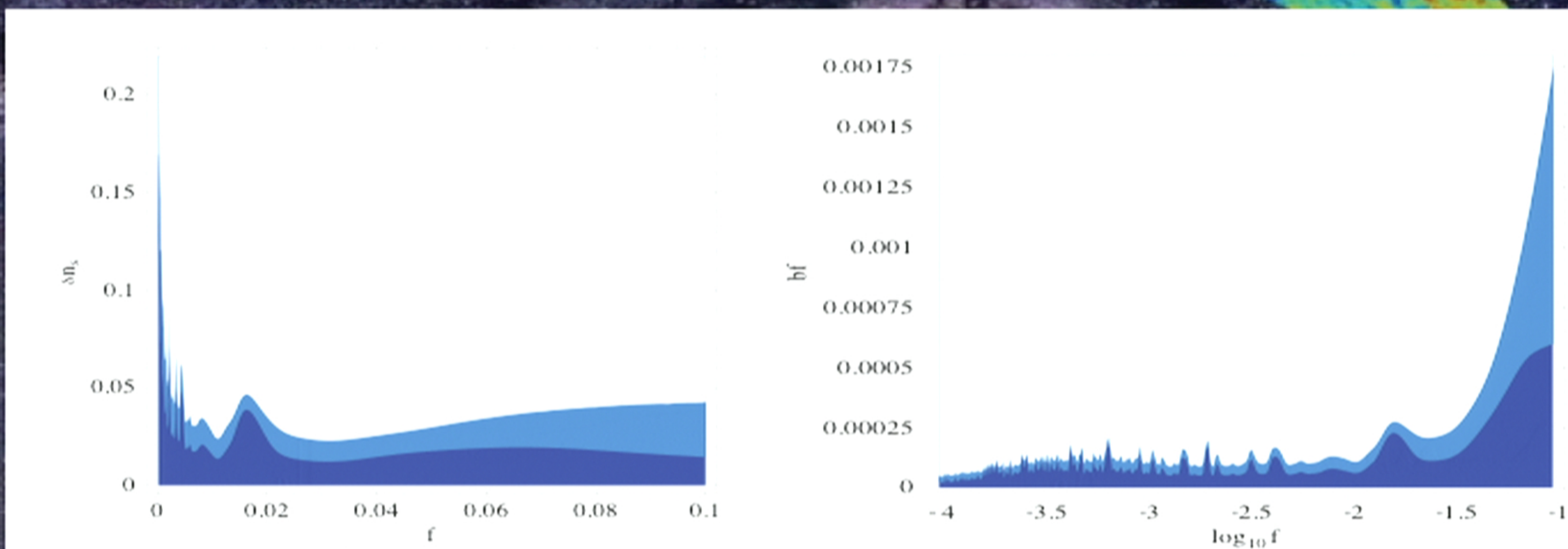
Constraints from WMAP5

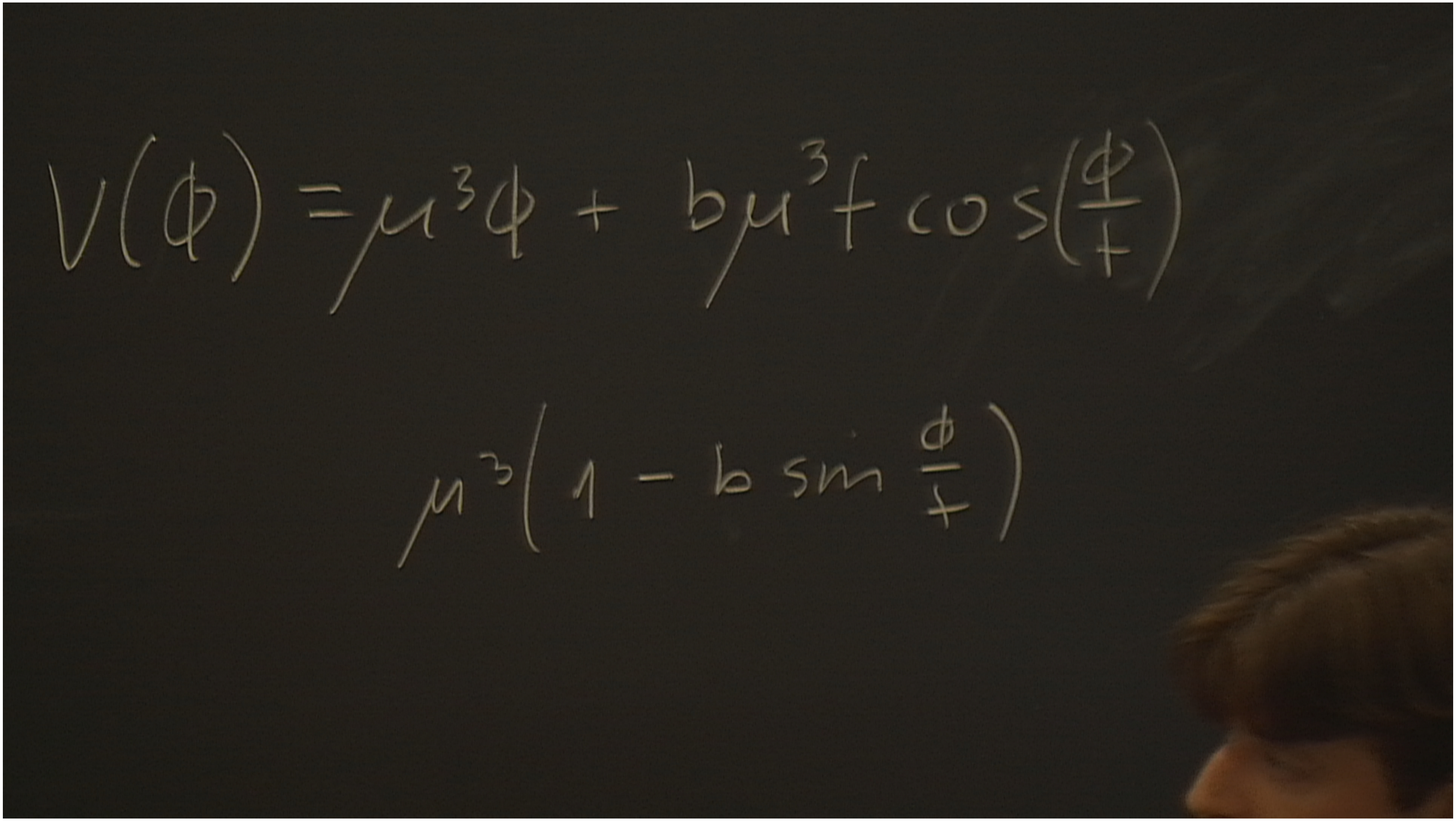



$V(\phi) = \mu^3\phi + b\mu^3 f \cos\left(\frac{\phi}{f}\right)$

Model Dependent Signatures I

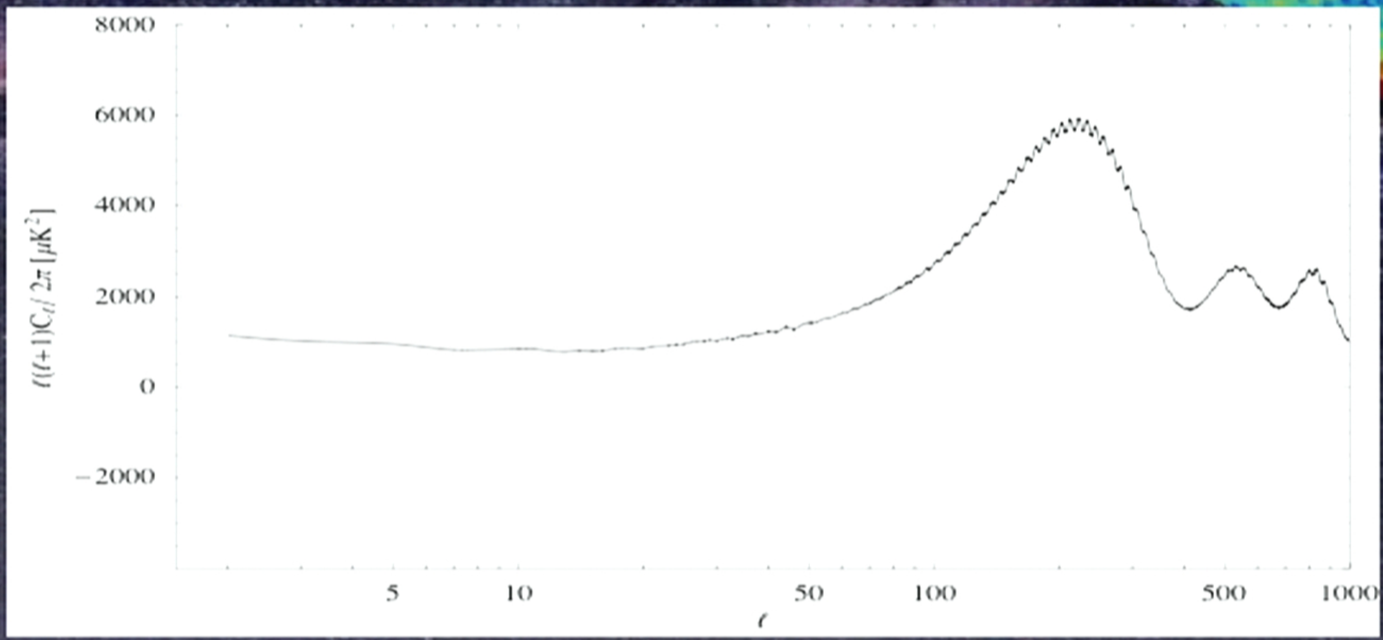
Constraints from WMAP5




$V(\phi) = \mu^3\phi + b\mu^3 f \cos\left(\frac{\phi}{f}\right)$
$\mu^3\left(1 - b \sin\frac{\phi}{f}\right)$

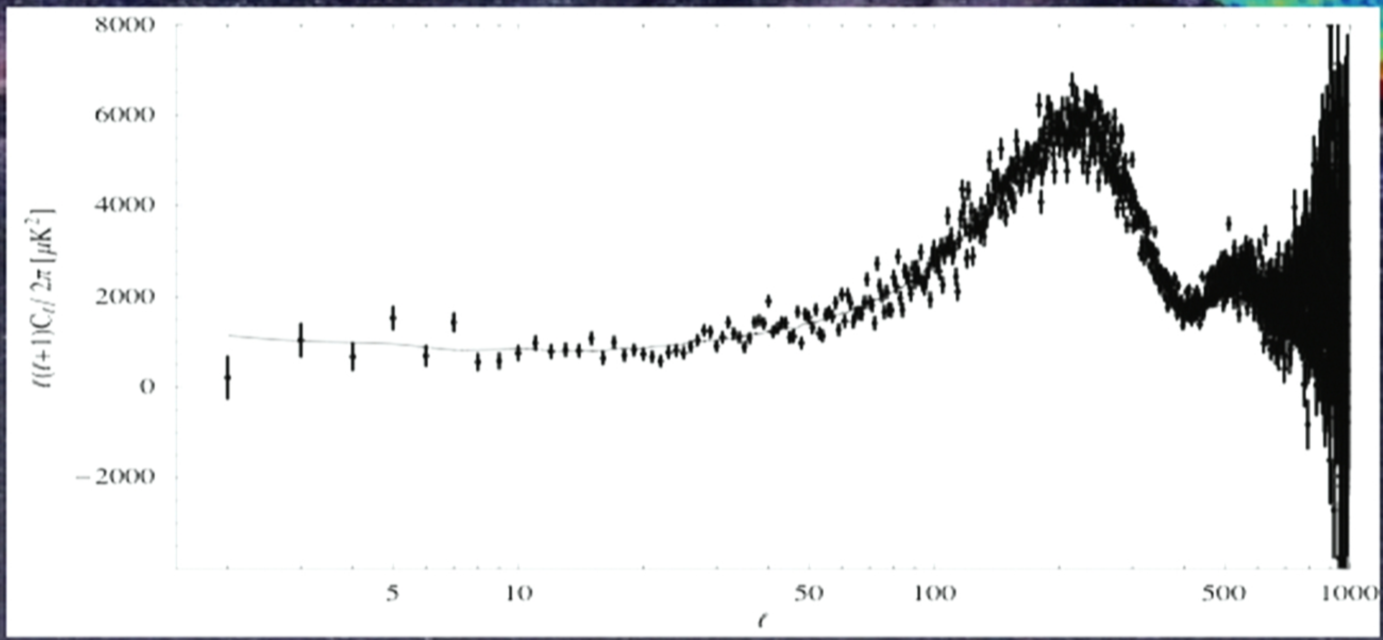
Model Dependent Signatures I

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

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Model Dependent Signatures II

Observable III: Resonant Non-Gaussianity

After some algebra

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

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

$$\mu^3\left(1 - b \sin \frac{\phi}{f}\right)$$

$$\langle \mathcal{R}(\vec{k}_1)\,\mathcal{R}(\vec{k}_2)\,\mathcal{R}(\vec{k}_3)\rangle = (2\pi)^7 \Delta_{\mathcal{R}}^4\, \delta(\vec{k}_1 + \vec{k}_2 + \vec{k}_3)$$

$$\frac{1}{k_1^2 k_2^2 k_3^2} \qquad \frac{\mathcal{G}(k_1, k_2, k_3)}{k_1 k_2 k_3}$$

Model Dependent Signatures II

Observable III: Resonant Non-Gaussianity

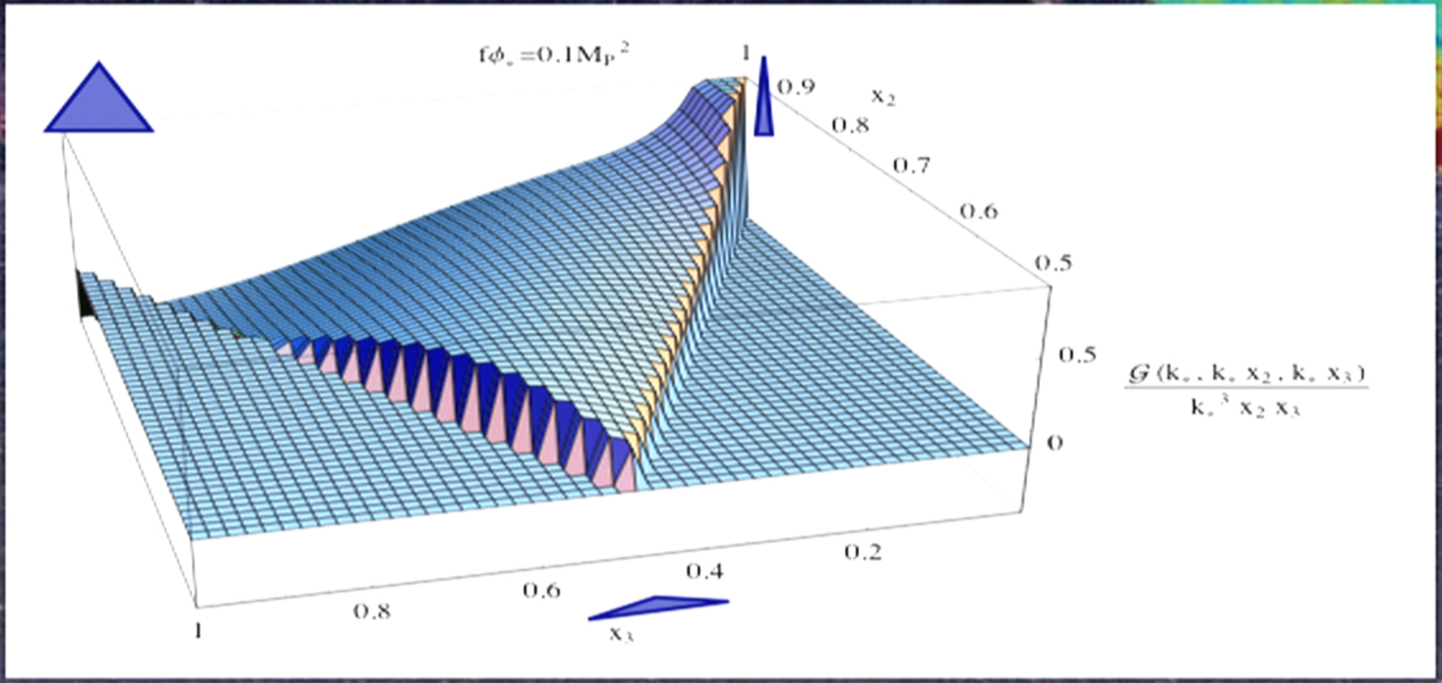
$$\frac{\mathcal{G}(k_1,k_2,k_3)}{k_1k_2k_3} = f^{\rm res}\left[\sin\left(\frac{\ln K/k_*}{f\phi_*}\right) + f\phi_*\sum_{i\neq j}\frac{k_i}{k_j}\cos\left(\frac{\ln K/k_*}{f\phi_*}\right)\right]$$

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

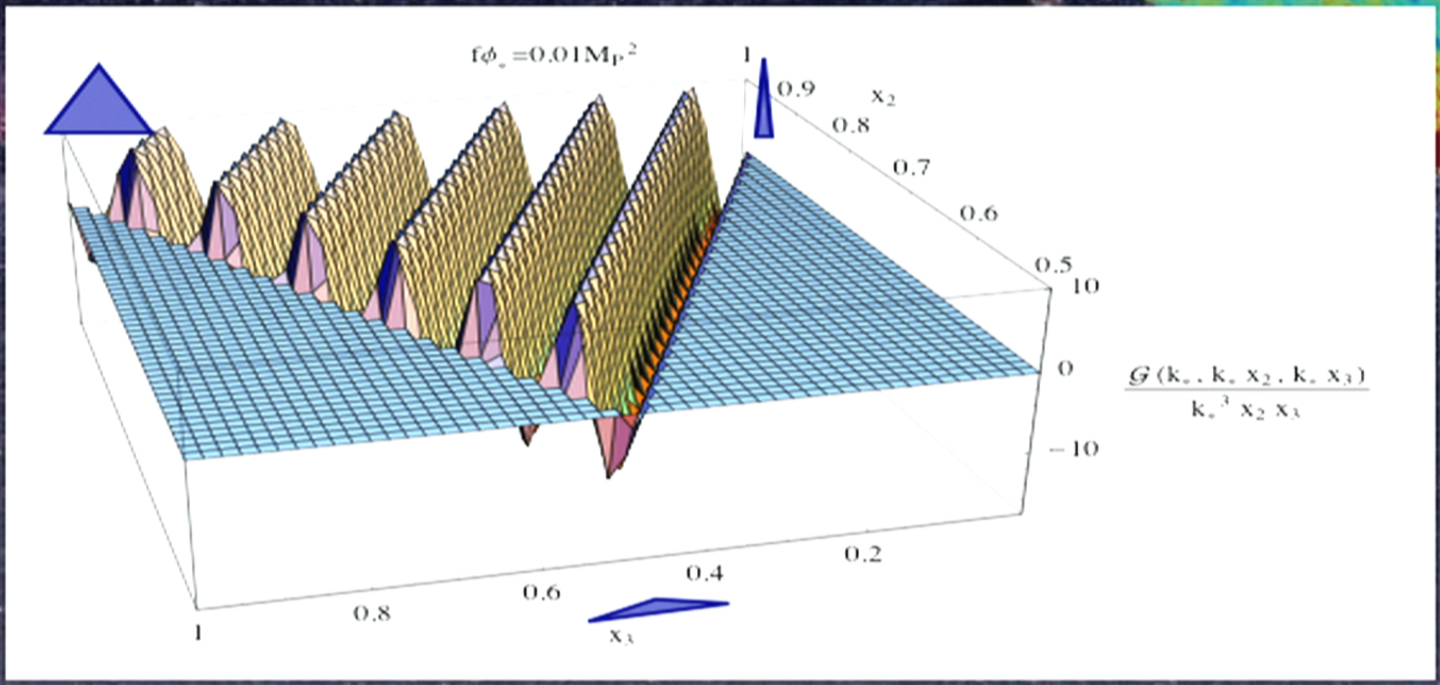
$$f^{\rm res} = \frac{3\sqrt{2\pi}b}{8(f\phi_*)^{3/2}}$$

(This satisfies the consistency condition.)

Model Dependent Signatures II



Model Dependent Signatures II



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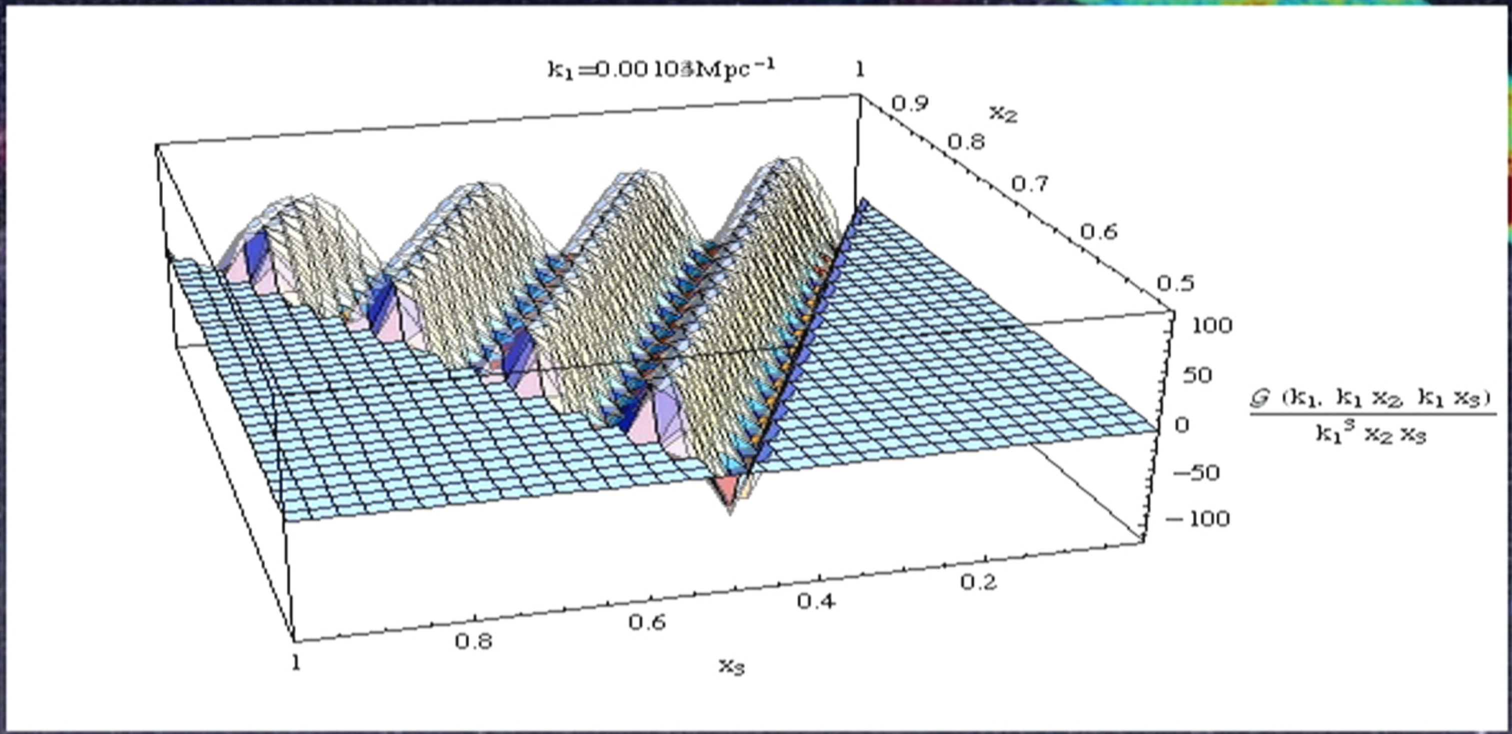

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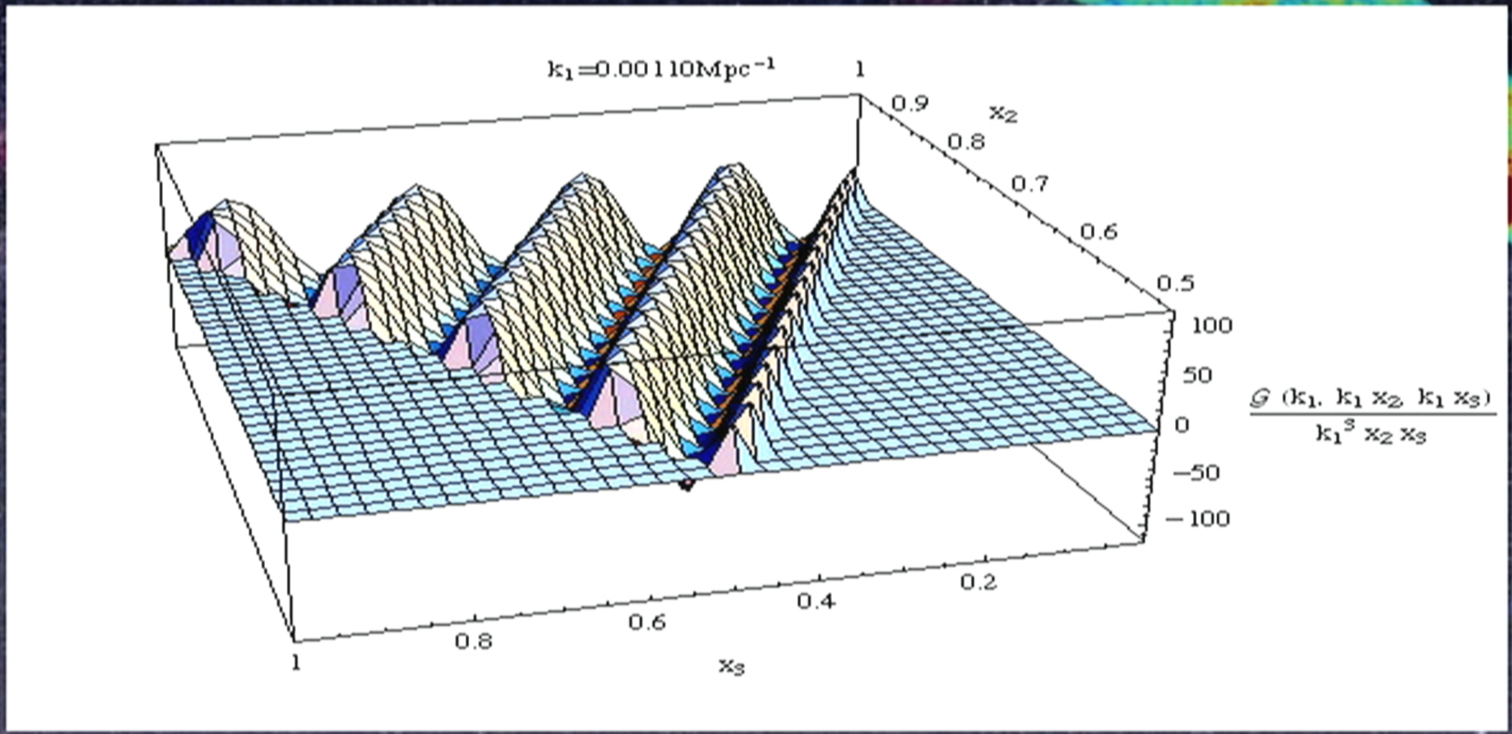

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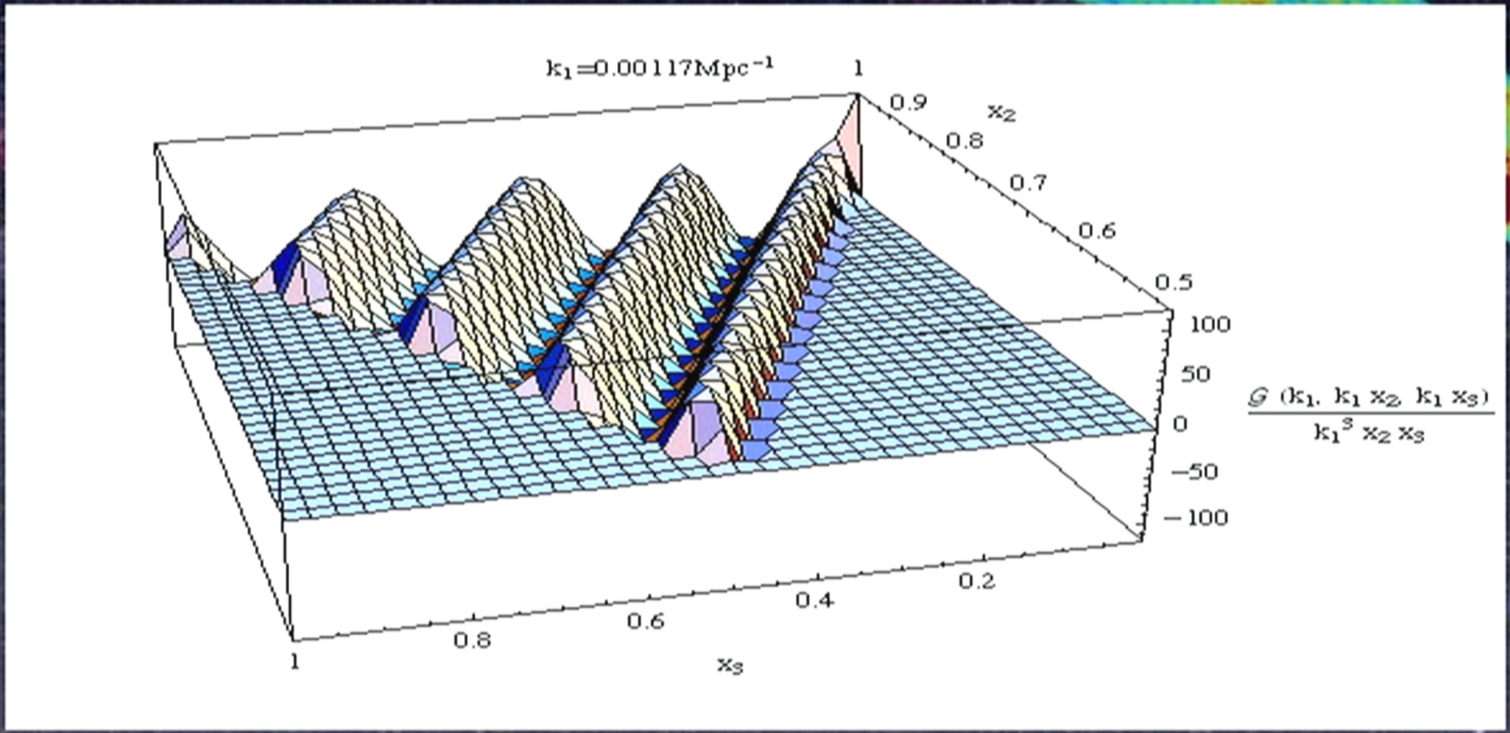

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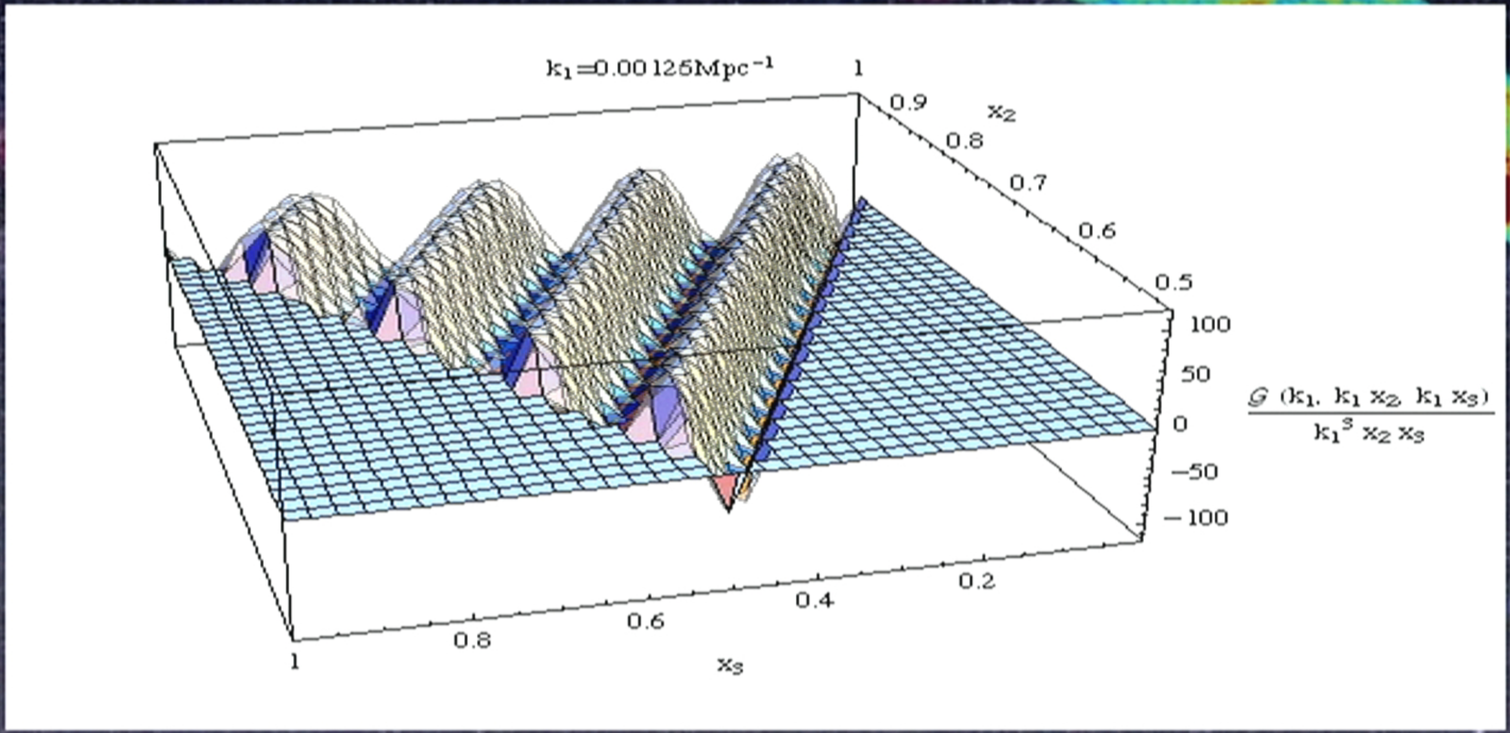

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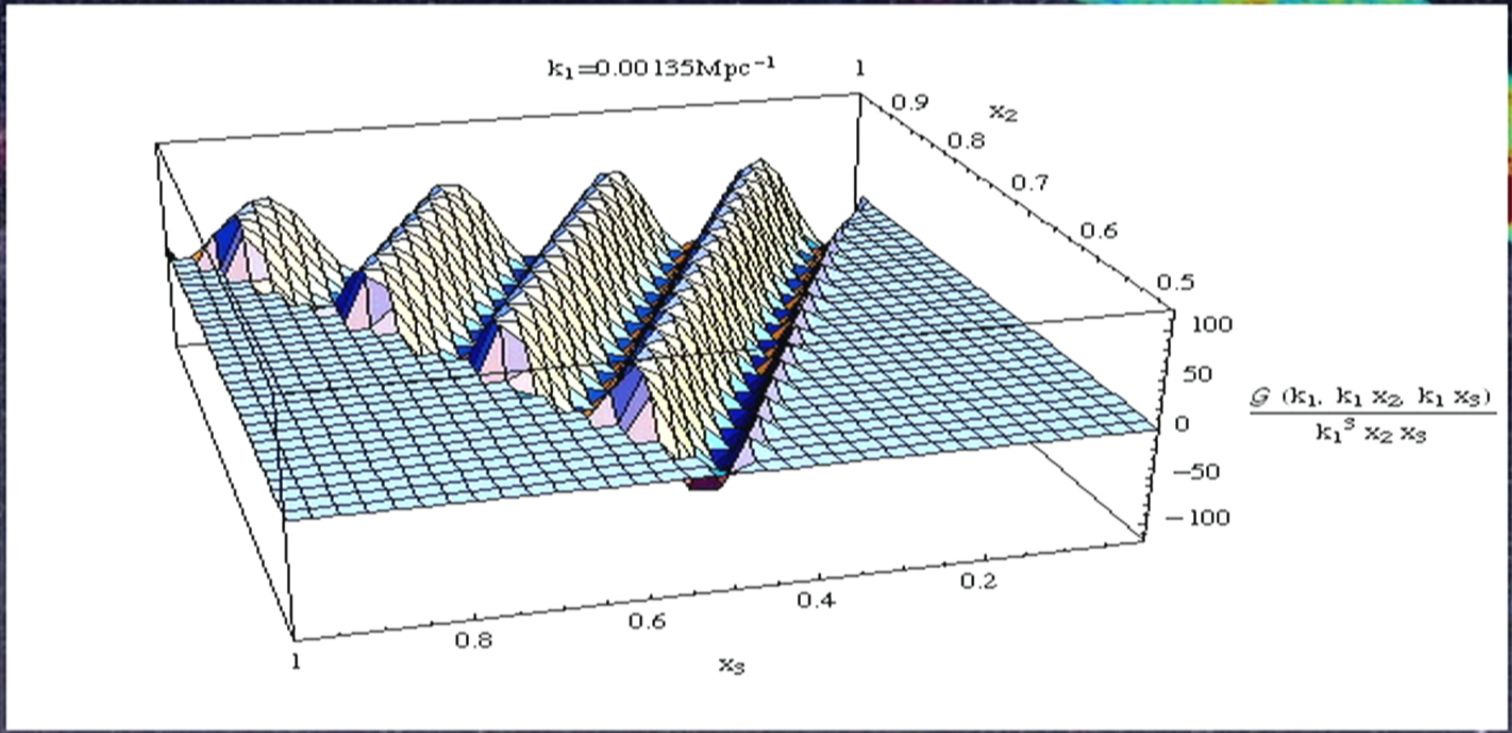

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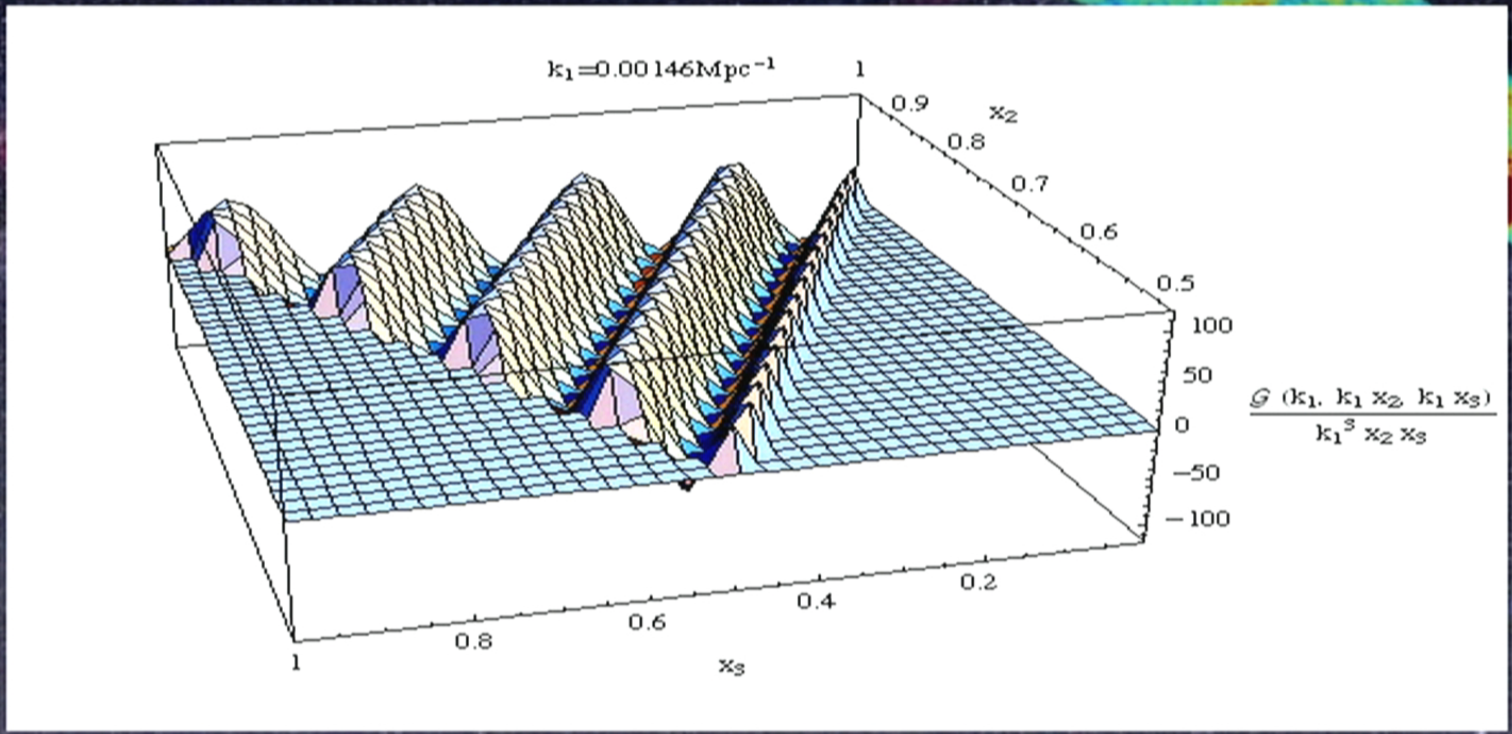

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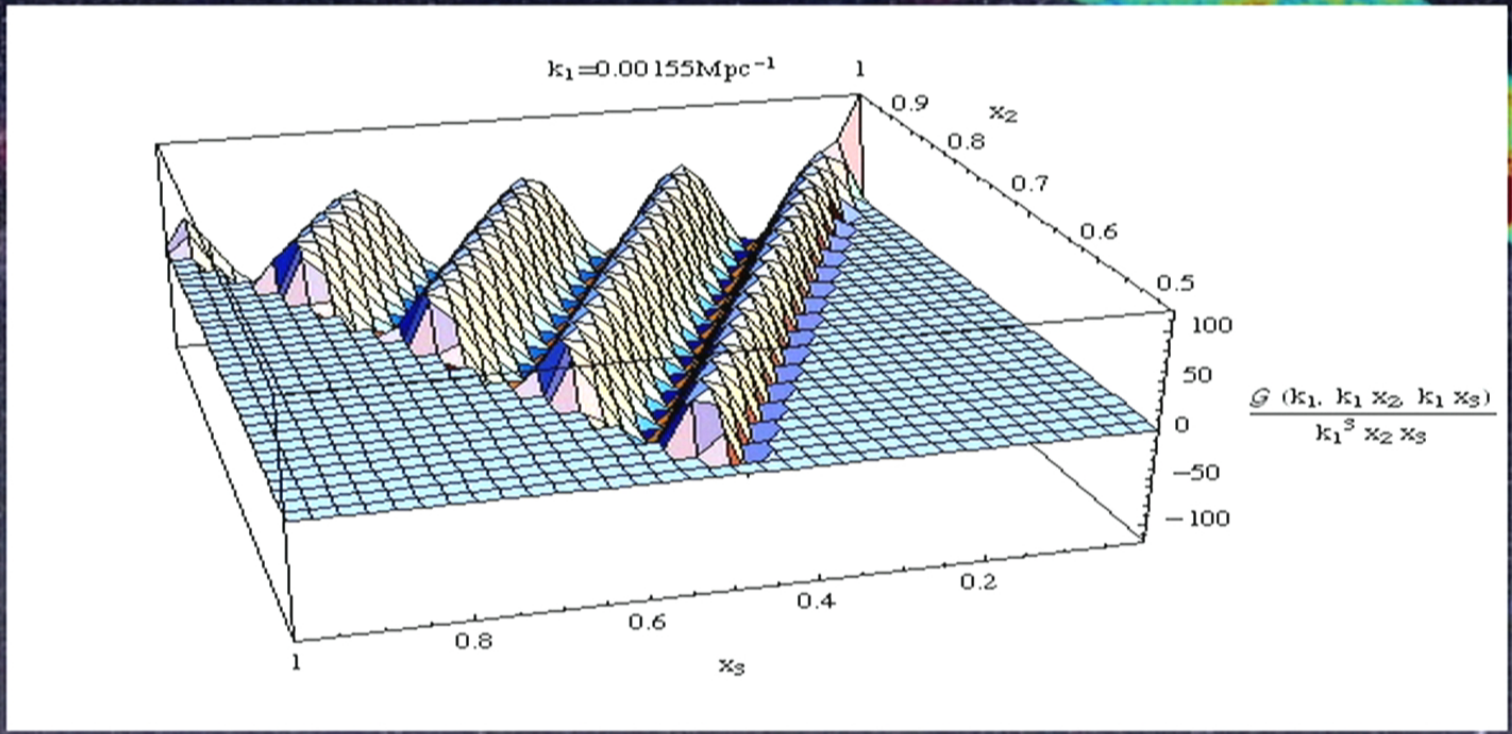

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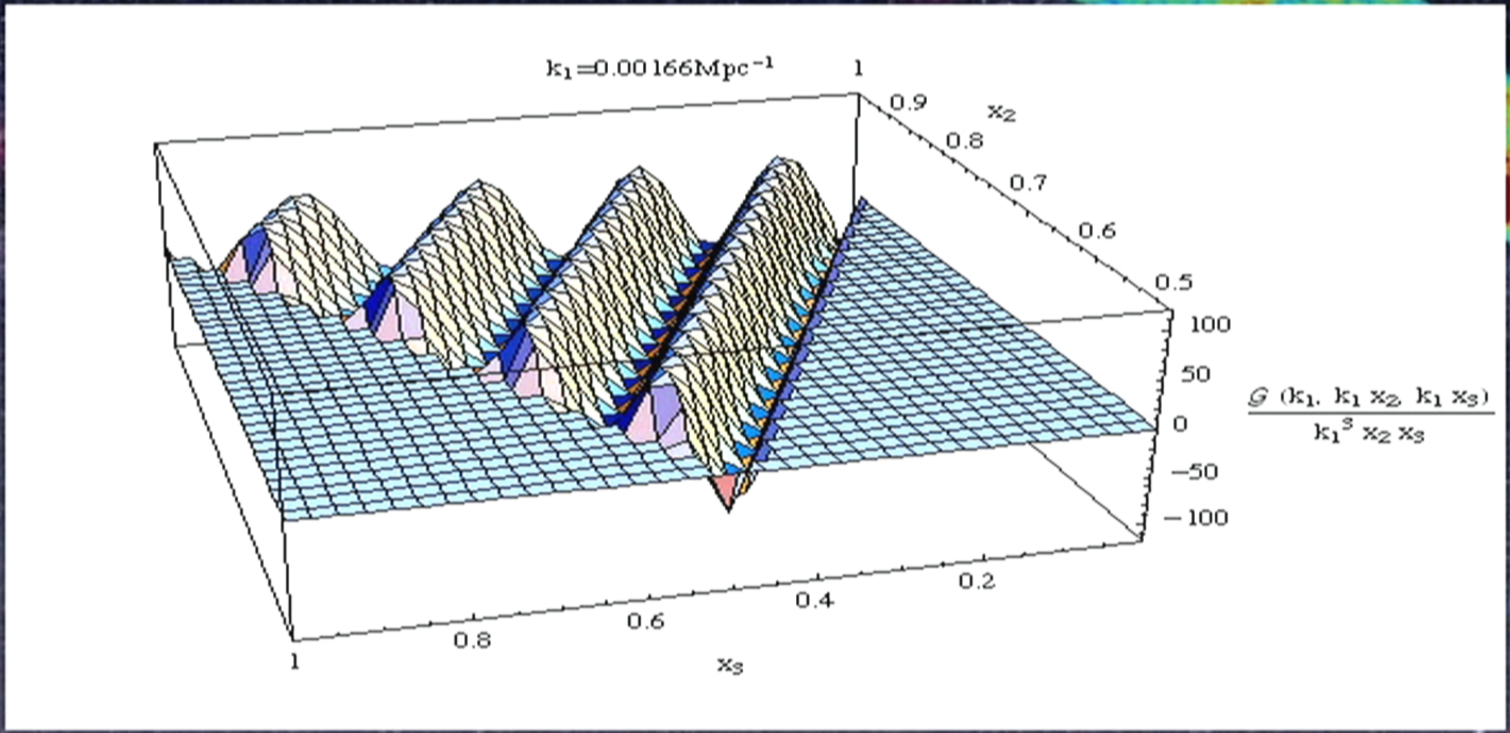

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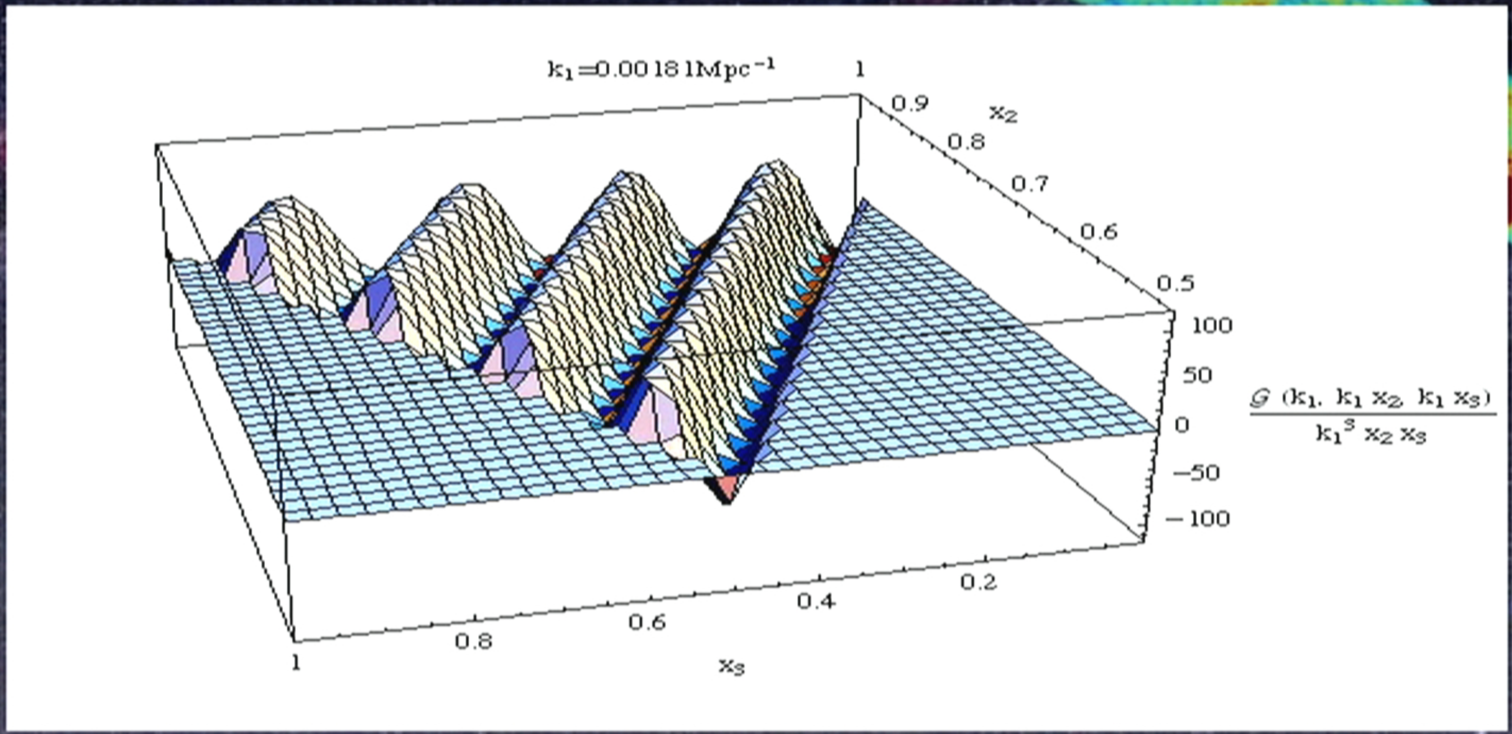

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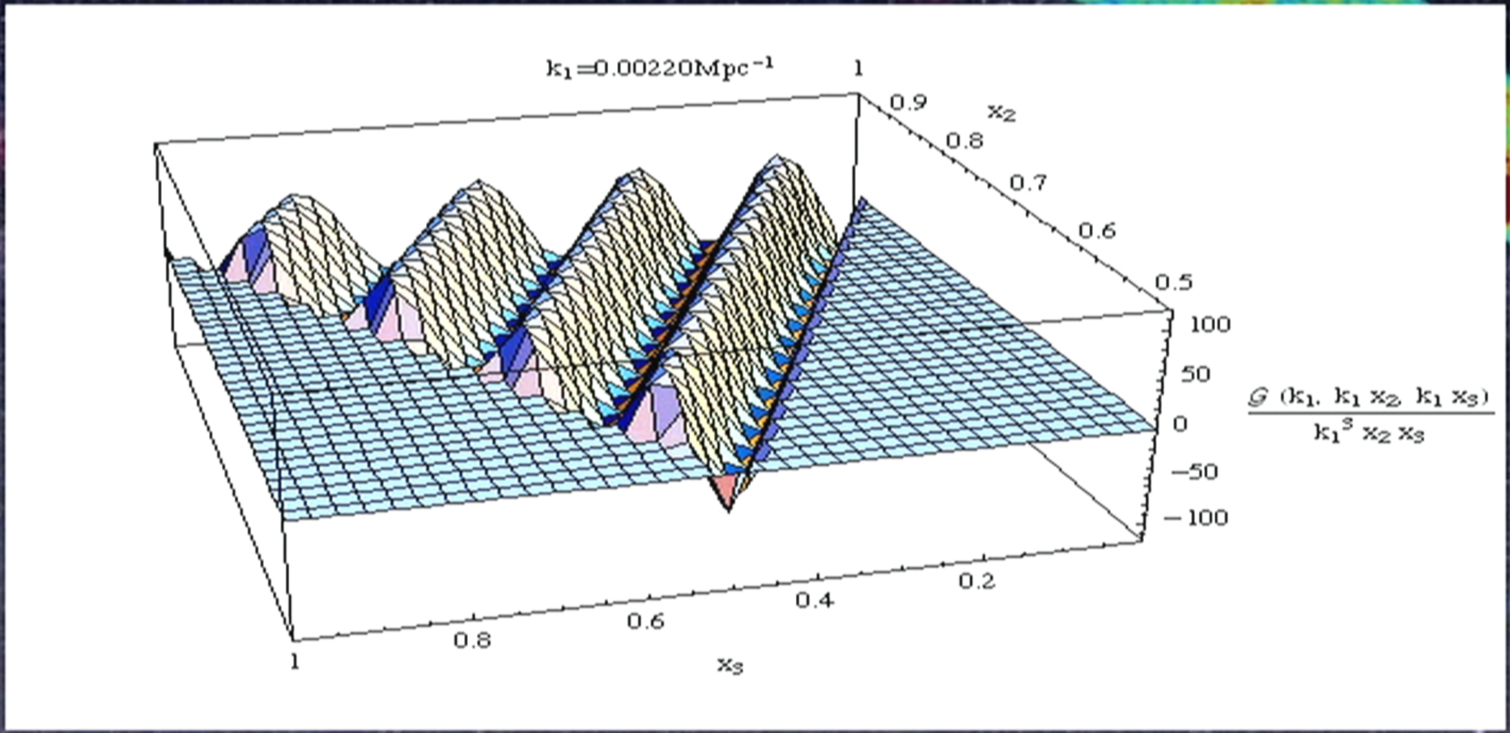
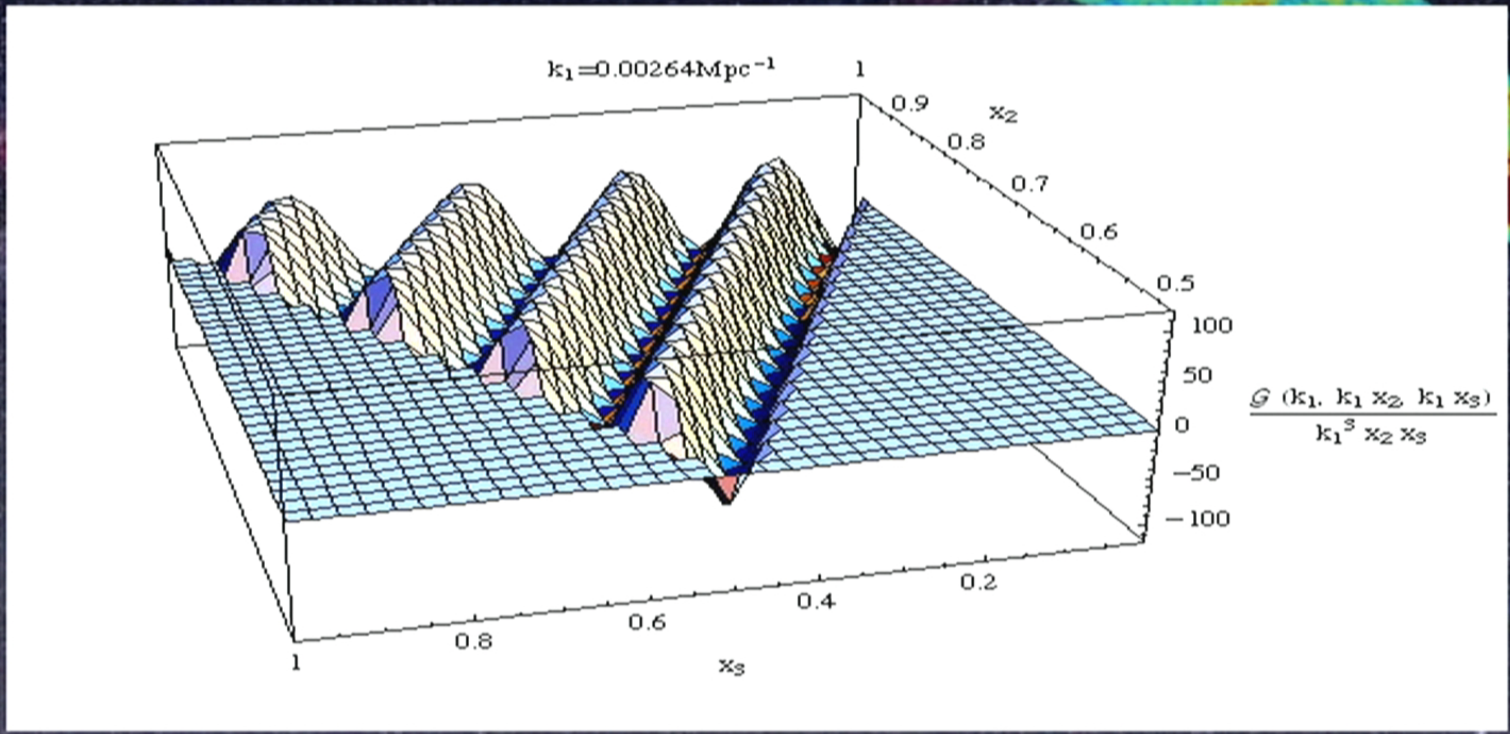
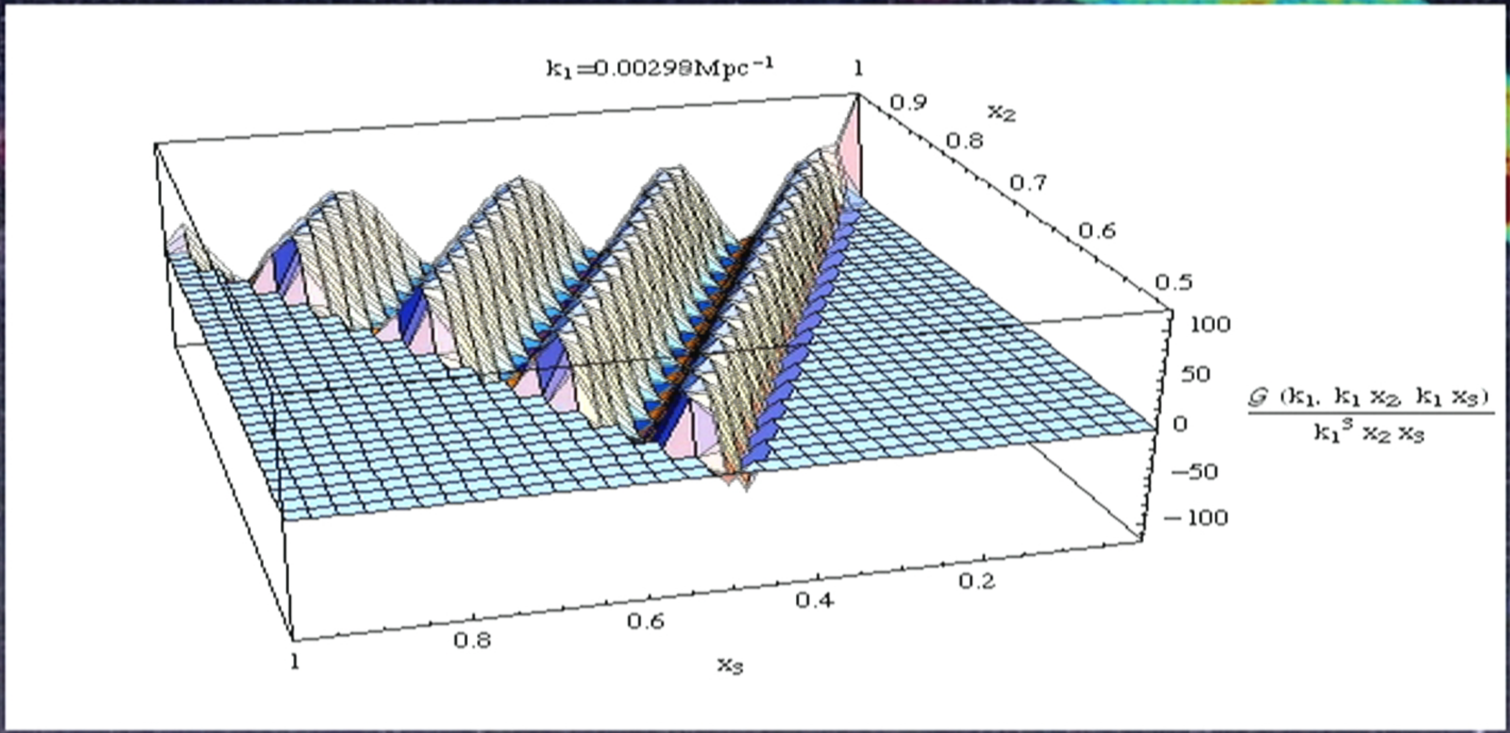


Model Dependent Signatures II

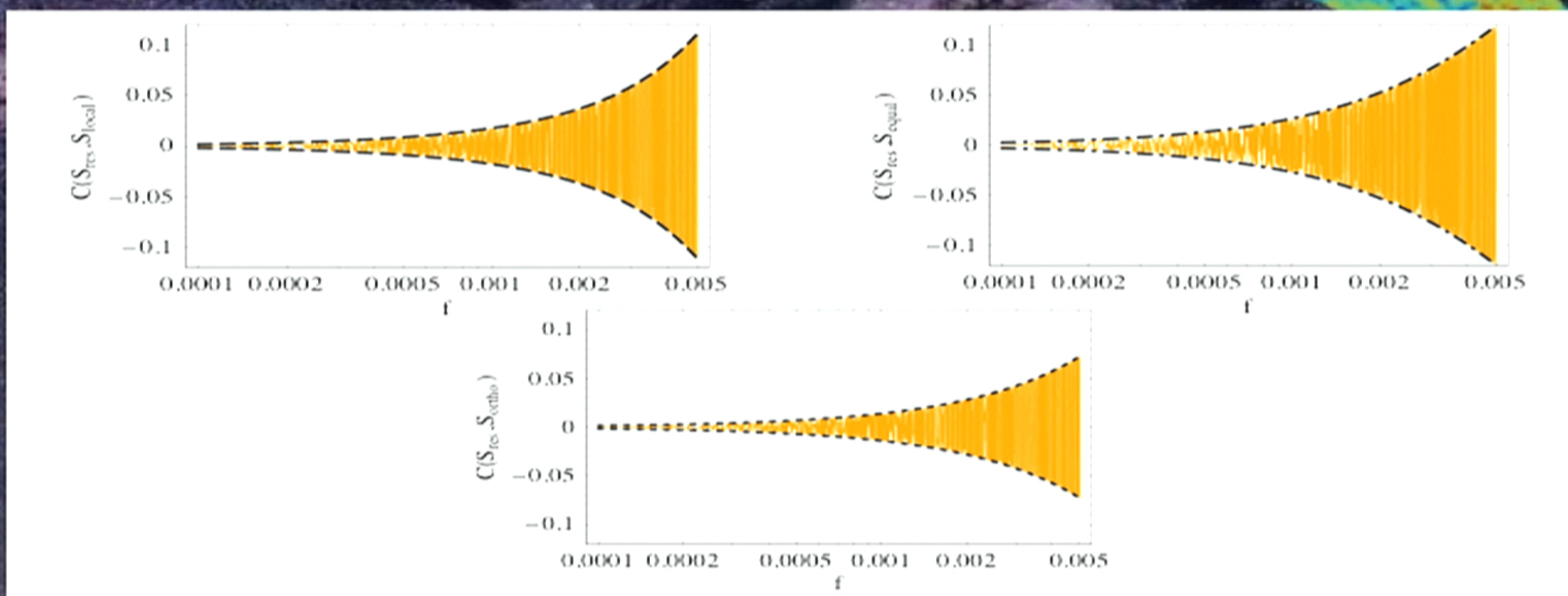


Model Dependent Signatures II



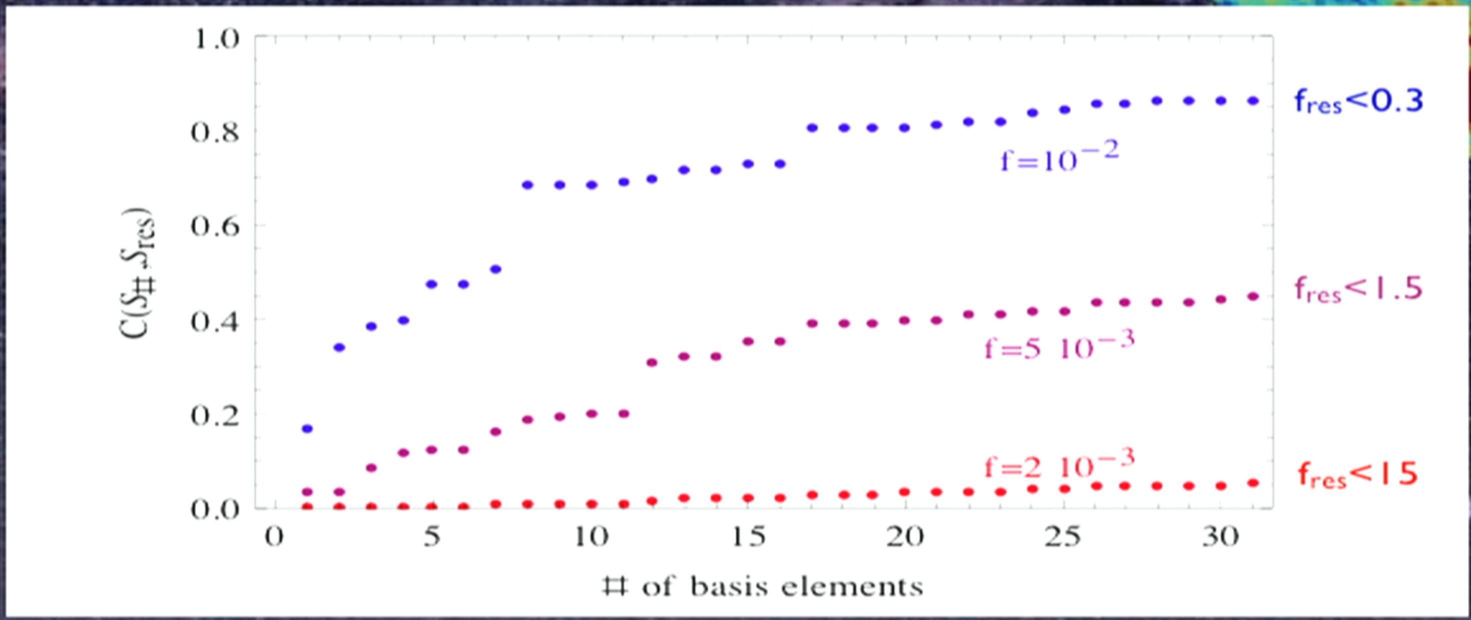
Model Dependent Signatures II

Can we convert existing constraints on local, equilateral, and orthogonal shapes into constraints on this shape?



Model Dependent Signatures II

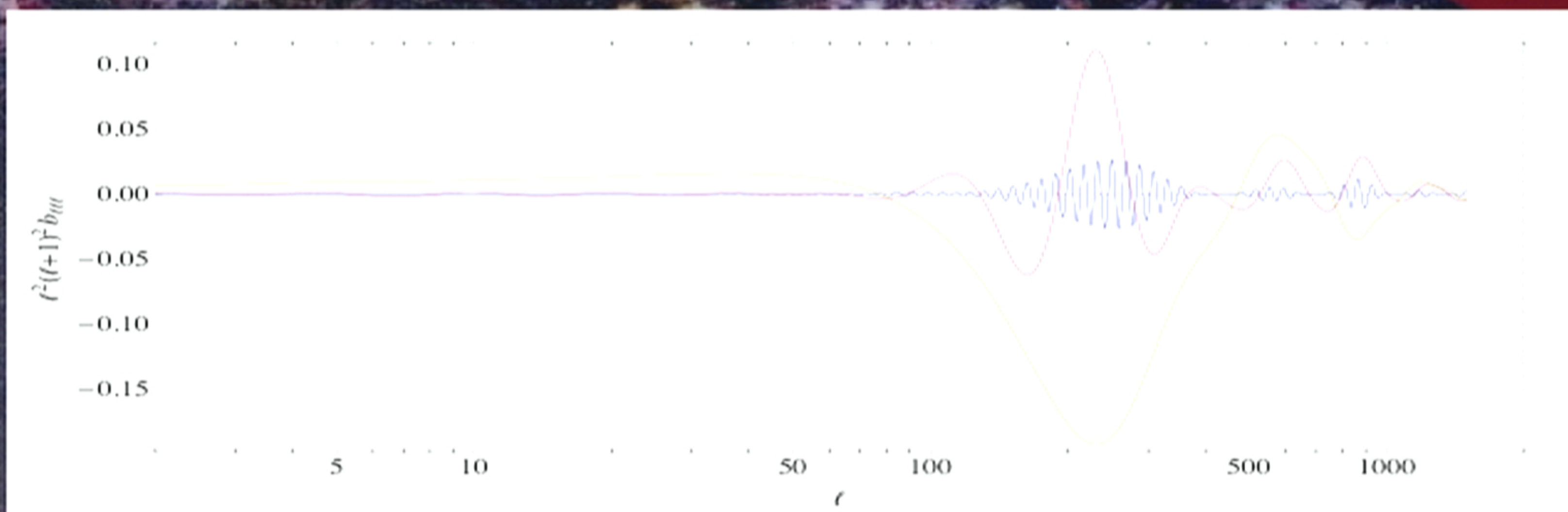
An expansion in a factorizable polynomial basis is limited to larger axion decay constants.

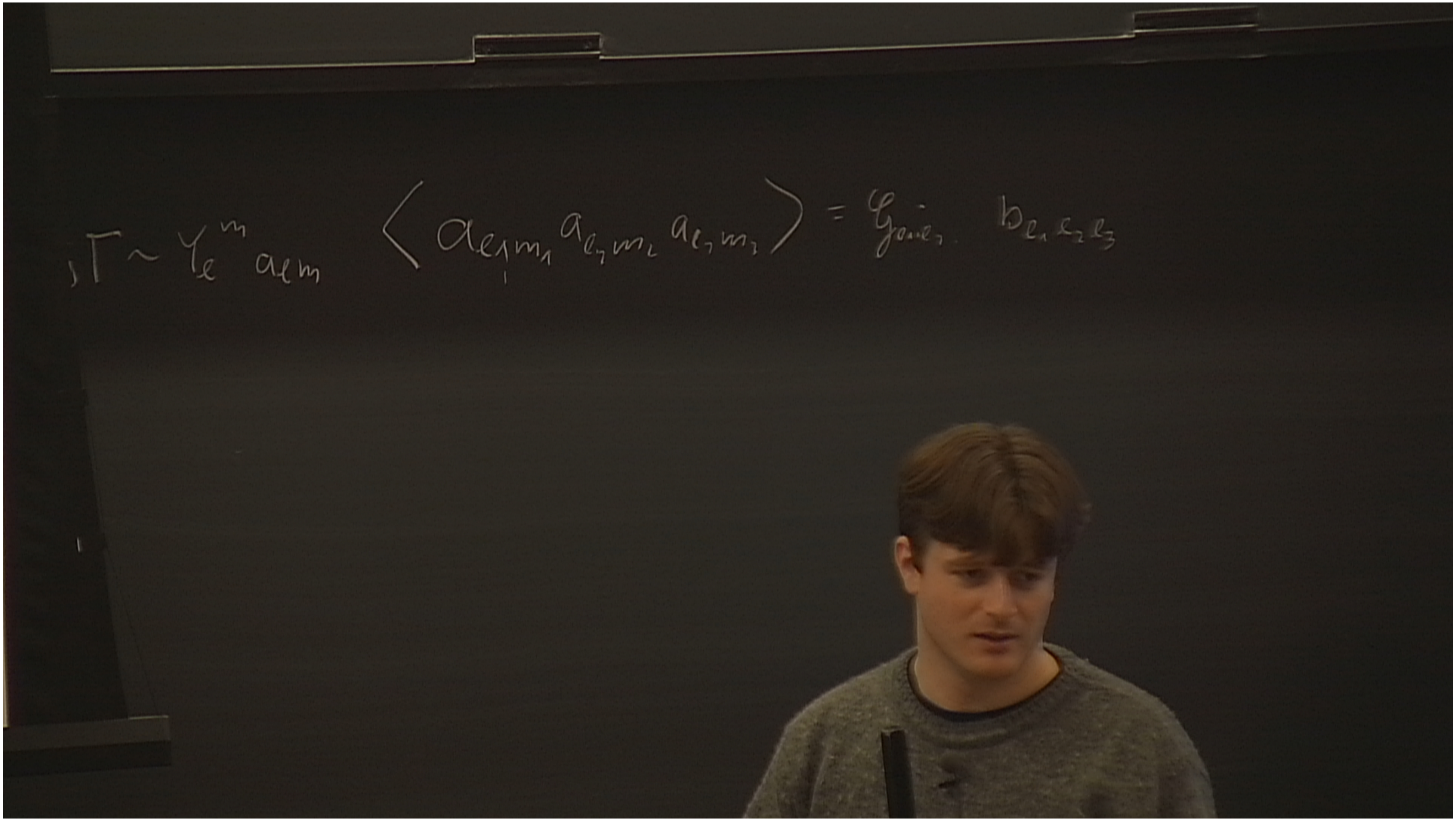


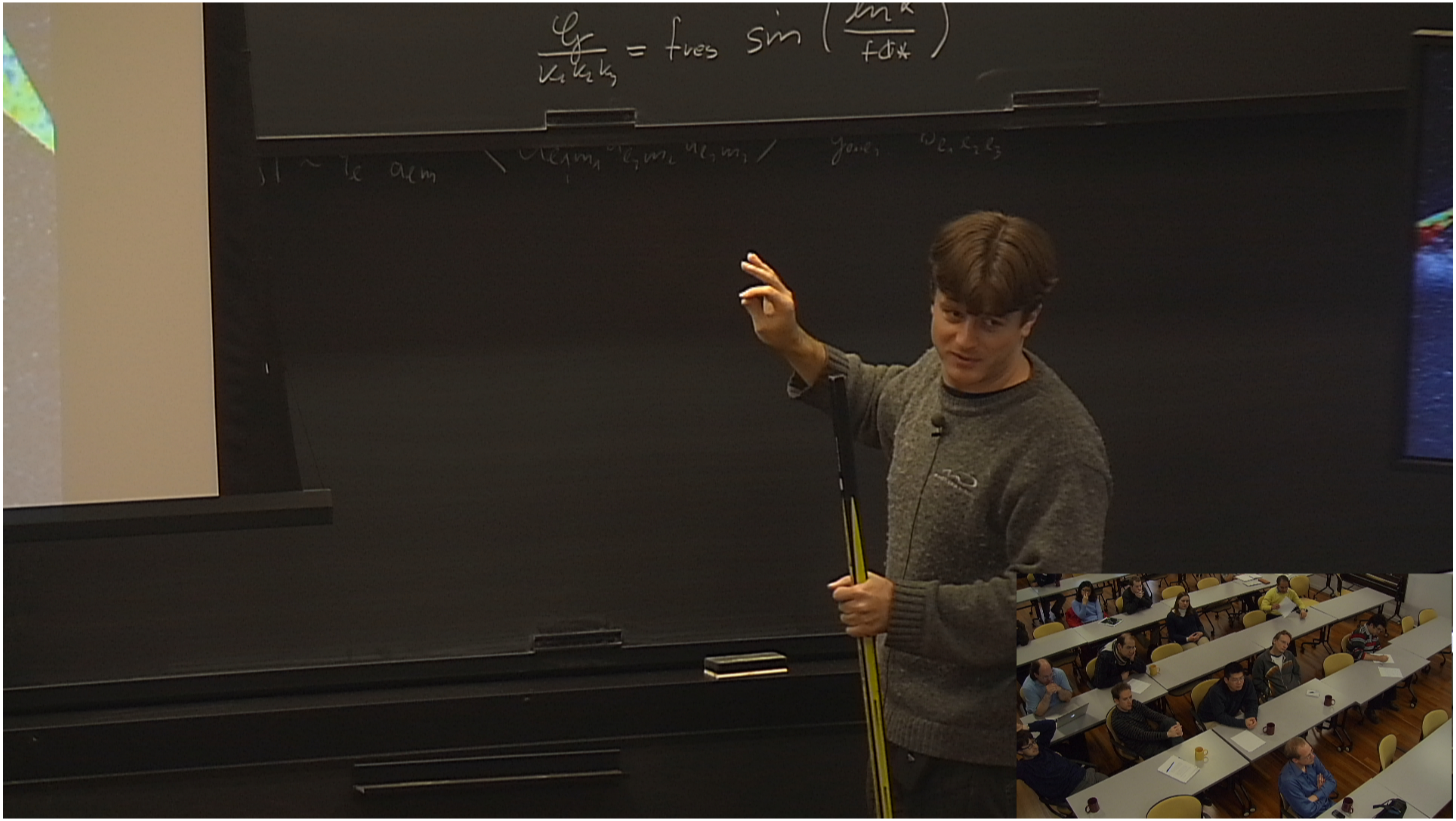
Model Dependent Signatures II

If one takes the best-fit point of the 2-pt analysis seriously, one expects a signal with f_{res}~1200.

Whether this is detectable remains to be seen.




Γ ~ Y_l^m a_lm
⟨a_{l1m1} a_{l2m2} a_{l3m3}⟩ = G b_{l1l2l3}



$$V(\phi) = \mu^3\phi + b\mu^3 f\cos\left(\frac{\phi}{f}\right)$$

$$\mu^3\left(1 - b\sin\frac{\phi}{f}\right)$$

$$\langle \mathcal{R}(\vec{k}_1)\,\mathcal{R}(\vec{k}_2)\,\mathcal{R}(\vec{k}_3)\rangle = (2\pi)^7\,\Delta_{\mathcal{R}}^4\,\delta(\vec{k}_1+\vec{k}_2+\vec{k}_3)$$

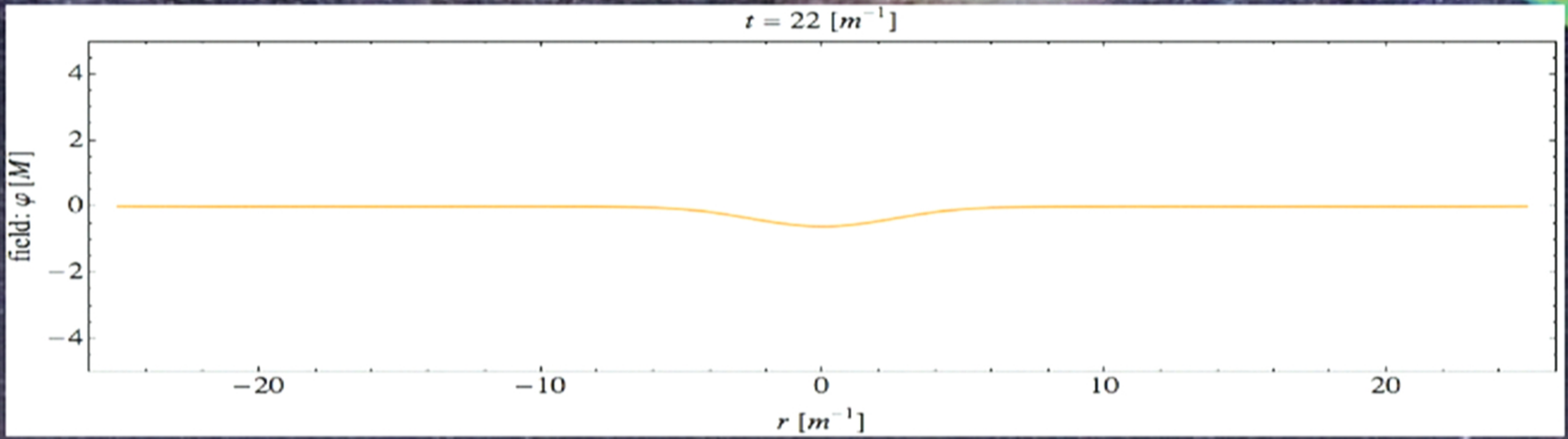
$$\frac{1}{k_1^2k_2^2k_3^2} \qquad \frac{\mathcal{G}(k_1,k_2,k_3)}{k_1\,k_2\,k_3}$$

$$\frac{\mathcal{G}}{k_1k_2k_3} = f_{res}\,\sin\left(\frac{\ln k}{f\phi_*}\right)$$

Model Dependent Signatures III

Oscillons at the end of inflation

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



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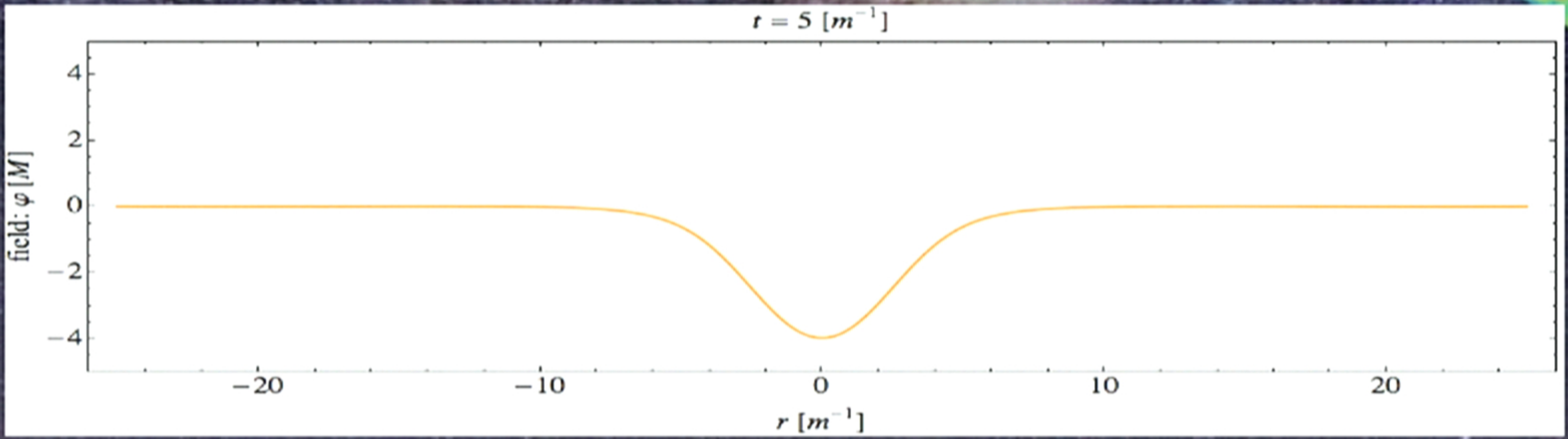

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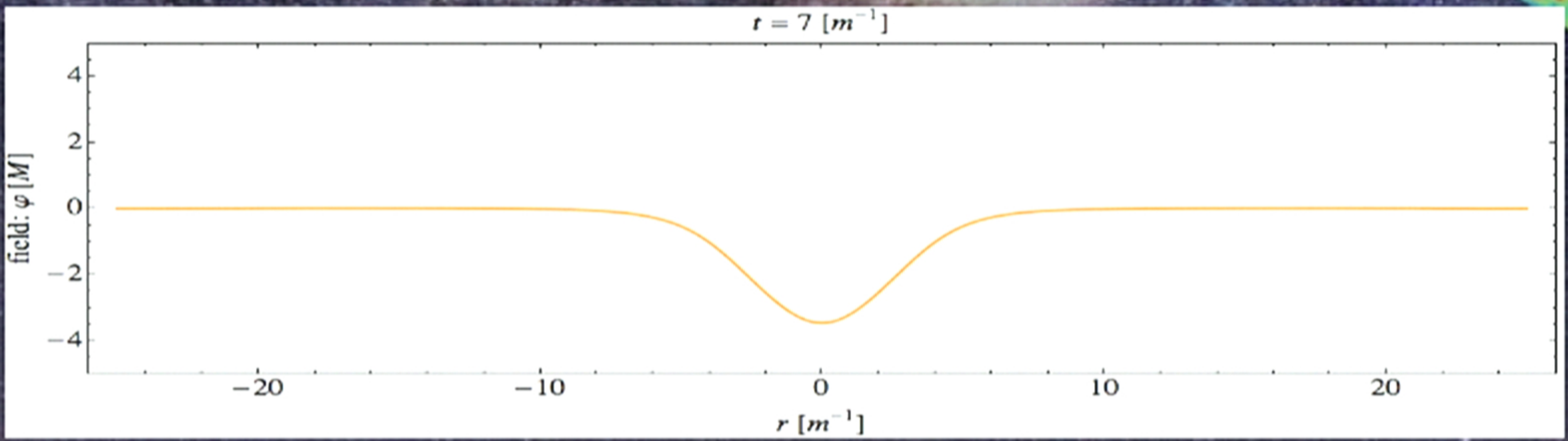
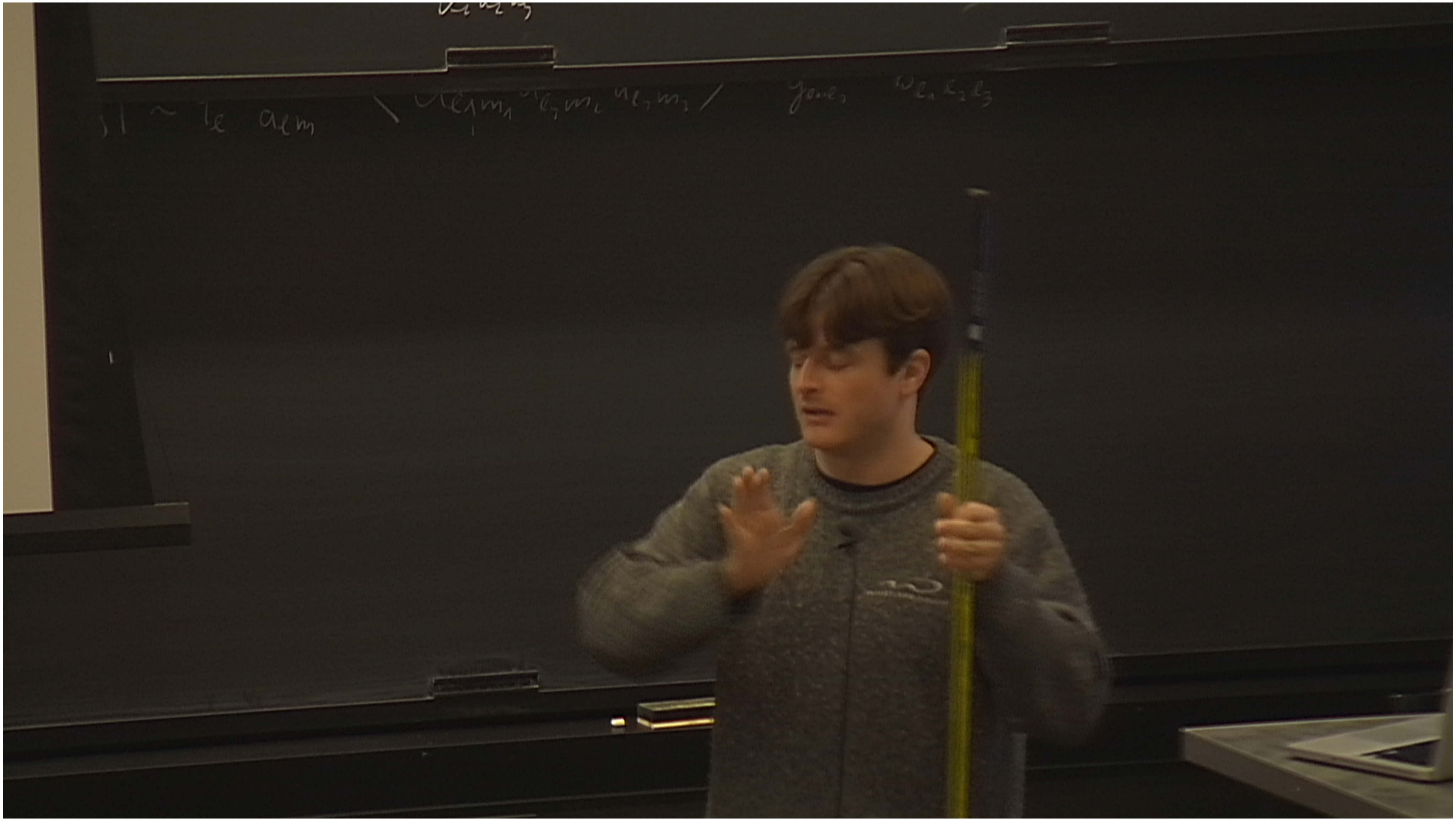




Model Dependent Signatures III

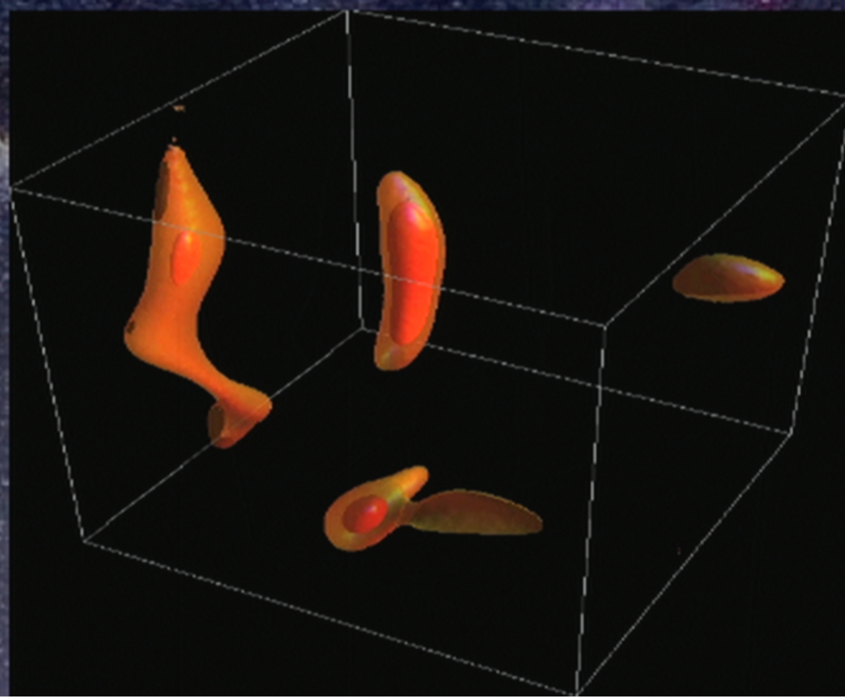
The calculations assume that the couplings to the degrees of freedom on the brane are small enough so that $\Gamma \ll H$ at the end of inflation.

Under these assumptions, oscillons are stable provided the potential flattens out for large field values.

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

Model Dependent Signatures III

Oscillon formation in the monodromy potential



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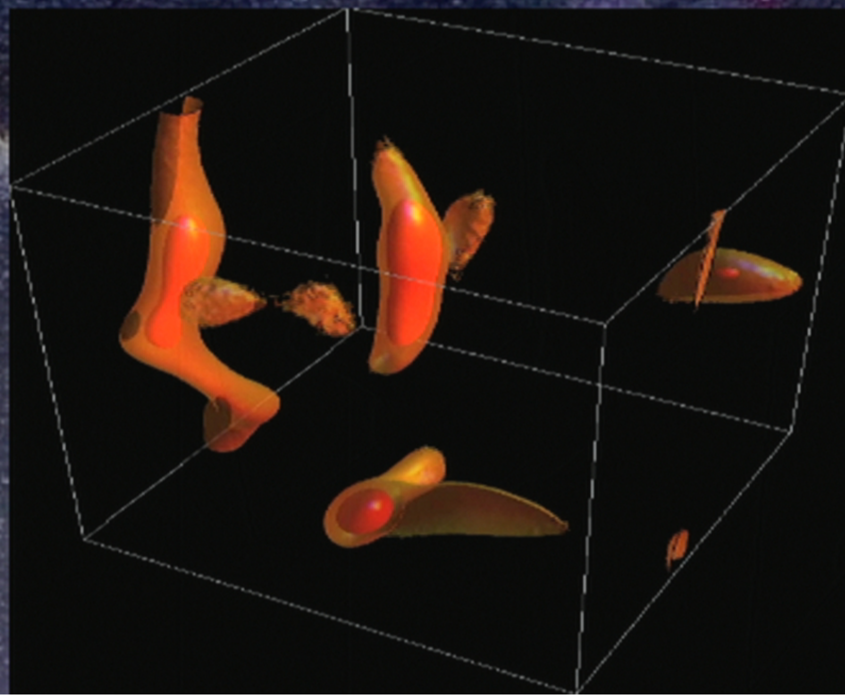
Model Dependent Signatures III

Oscillon formation in the monodromy potential



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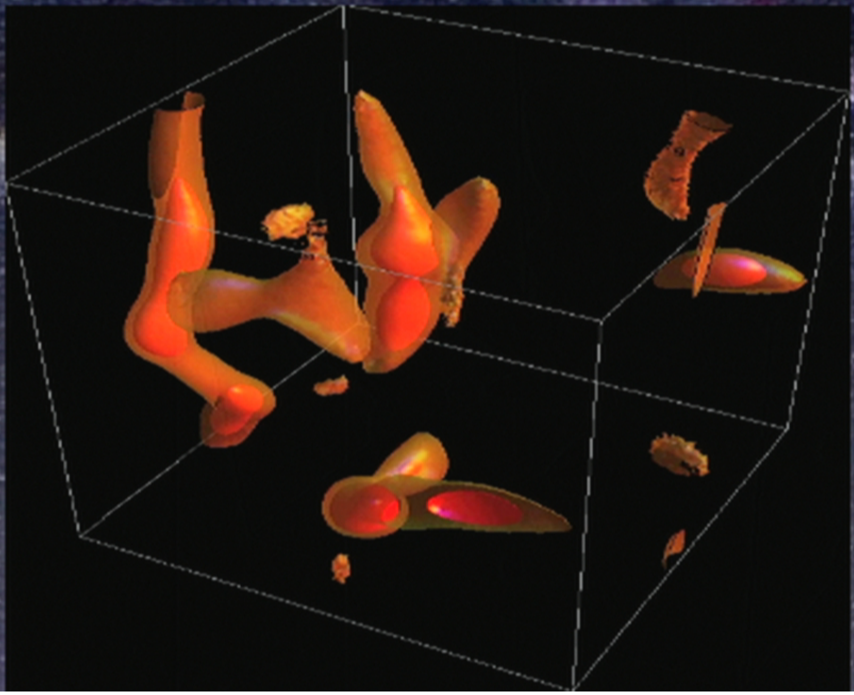
Model Dependent Signatures III

Oscillon formation in the monodromy potential



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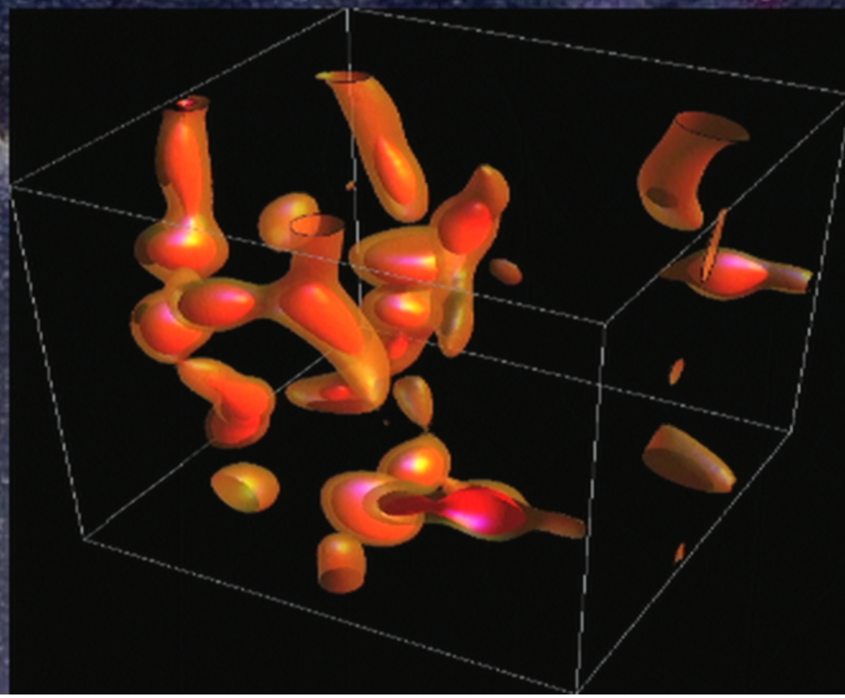
Model Dependent Signatures III

Oscillon formation in the monodromy potential



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Conclusions

- String theory seems to possess 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


Thank you