Title: A comment on gravitational waves and the scale of supersymmetry breaking

Date: Jun 16, 2010 03:45 PM

URL: http://pirsa.org/10060017

Abstract:
It has been suggested, by Kallosh and Linde, that a generic bound on inflation in string theory keeps the Hubble scale of inflation H smaller than the gravitino mass, $m_{3/2}$. Given that models with low-energy supersymmetry have $m_{3/2}$ smaller than a TeV, this is a severe constraint, and would suggest that one is forced to choose between high-scale inflation and low-scale supersymmetry. The bound arises by considering possible decompactification instabilities of the extra (compactified) dimensions of string theory, during the inflationary epoch. I explain the arguments that give rise to such a bound, and describe recent work with T. He and A. Westphal exhibiting large-field chaotic inflation models in string-inspired supergravities that have $H >> m_{3/2}$ but avoid decompactification. I conclude that even within the framework of string theory, high-scale inflation and low-energy supersymmetry may well be compatible.

Gravity Waves & the LHC:

Towards High-Scale Inflation with low-energy SUSY

Alexander Westphal
Stanford University
(arXiv: 1003.4265)

with: Temple He & Shamit Kachru

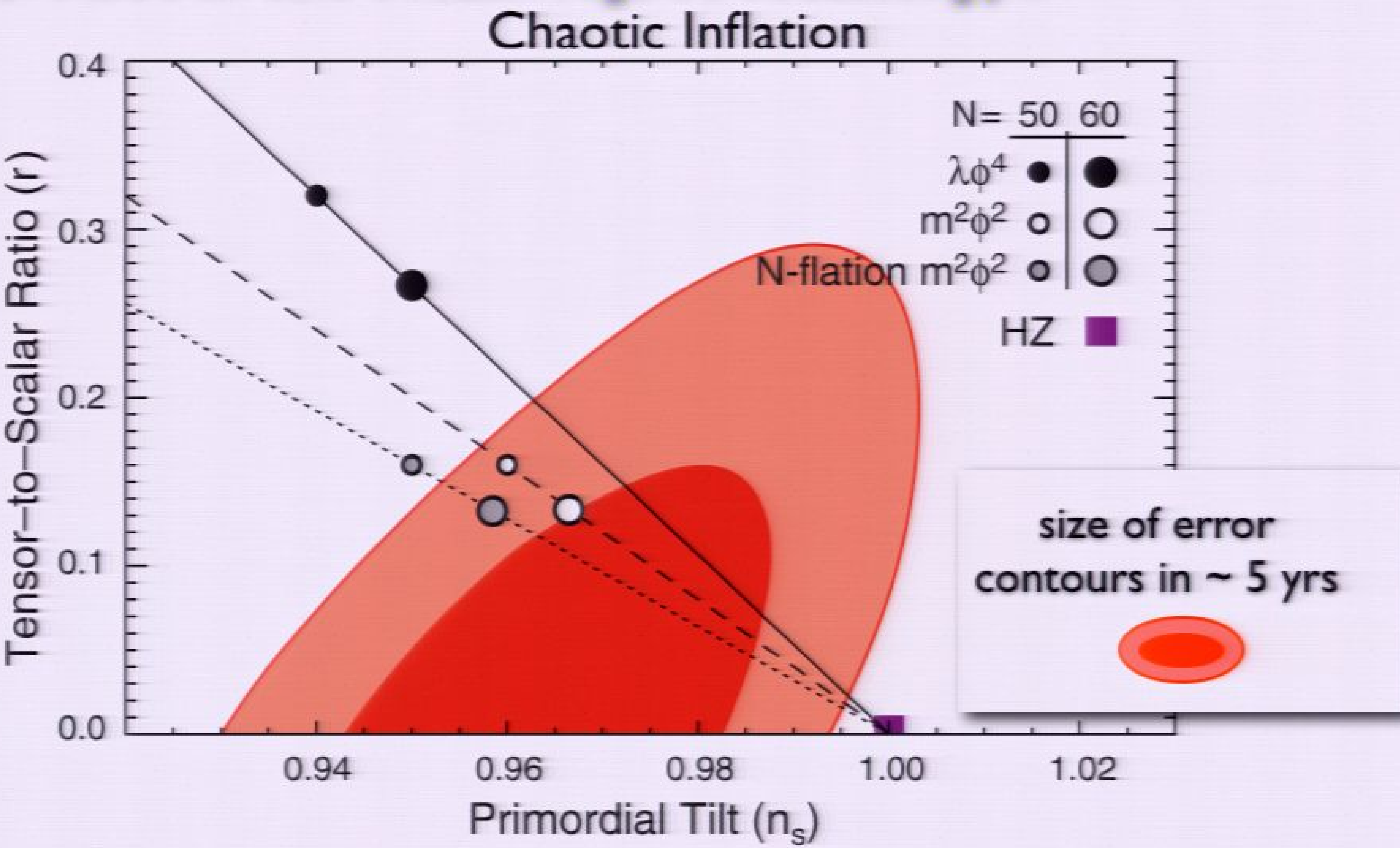
where I want to take you ...

- why:
  - large-field inflation (Φ moves more than M_P)?
  - strings?

- inflation & moduli stabilization - the Kallosh-Linde problem

- the demise of the problem - natural high-scale inflation @ the TeV
  - a natural setup for $H >> m_{3/2}$ in KKLT
  - dynamics of the volume modulus during inflation
  - hierarchies & scales - horse trading

We live in the Golden Age of cosmology !



expect dramatic improvement in next 5 yrs:

Planck & **BICEP2** taking data, Keck Array ('10..)

Inflation ...

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$$\Rightarrow \quad \epsilon \equiv -\frac{\dot{H}}{H^2} \simeq \frac{1}{2}\left(\frac{V'}{V}\right)^2 \ll 1 \quad , \quad \eta \equiv \frac{\dot{\epsilon}}{\epsilon H} \simeq \frac{V''}{V} \ll 1$$

with the Hubble parameter $H^2 = \frac{\dot{a}^2}{a^2} \simeq const. \sim V$

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- *but:* if field excursion sub-Planckian, no measurable gravity waves:

$$r \equiv \frac{\mathcal{P}_T}{\mathcal{P}_S} \leq 0.003 \left(\frac{50}{N_e}\right)^2 \left(\frac{\Delta\phi}{M_P}\right)^2$$

why strings?

- *large field* model of inflation, i.e. "chaotic inflation"

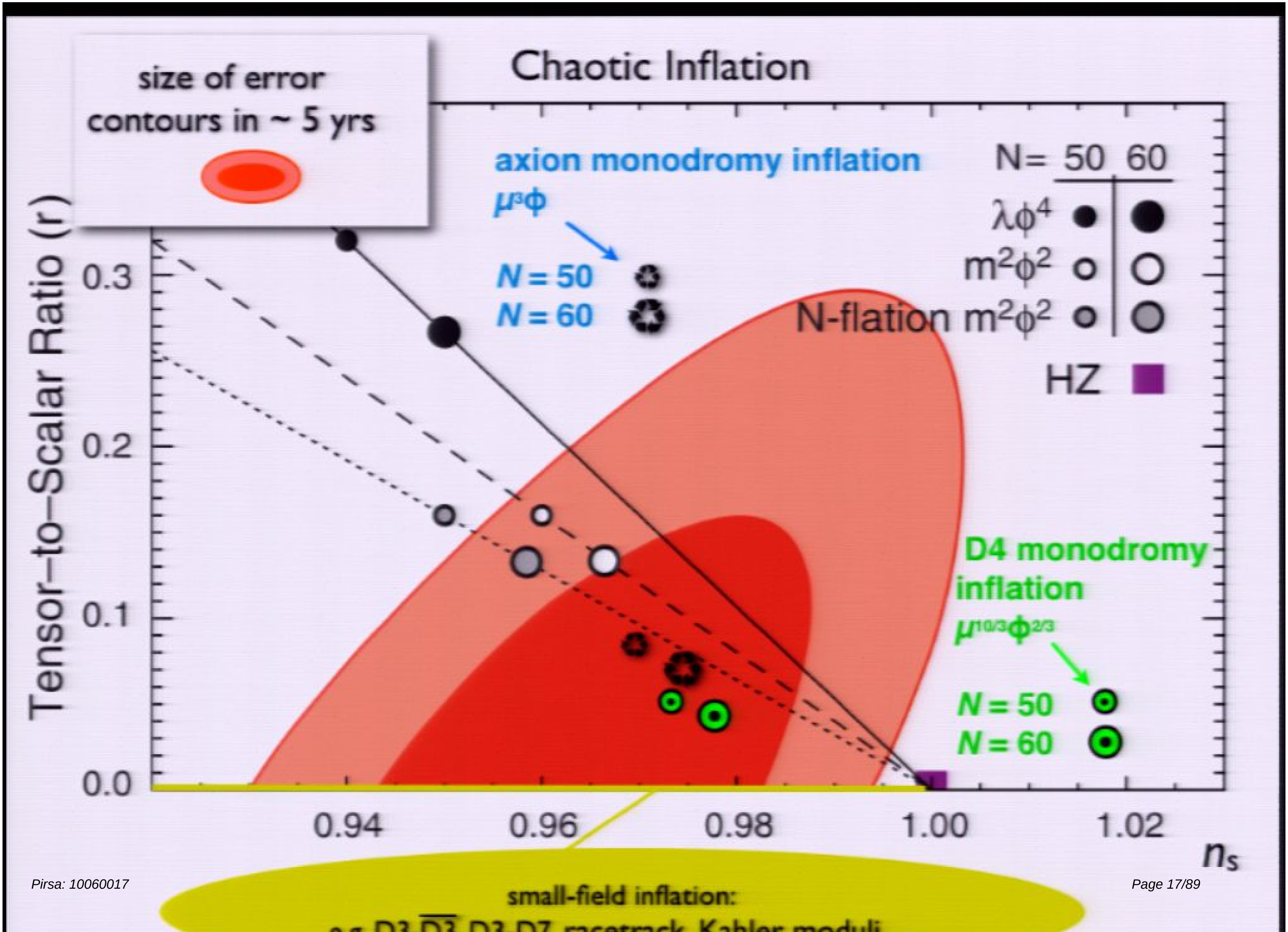
$$\Delta\phi > M_P \quad \Rightarrow \quad r > 0.01$$

- with control of ε & η over a super-Planckian field distance - avoid generic dim ≥ 6 operators:

$$\delta V \sim V(\phi)\frac{(\phi-\phi_0)^2}{M_{\rm P}^2}$$

need UV-complete theory: e.g. strings

- idea: arrange for approximate shift symmetry of ϕ , broken only by the inflaton potential itself
[Linde '83]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μ³φ
N = 50
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N= 50 60
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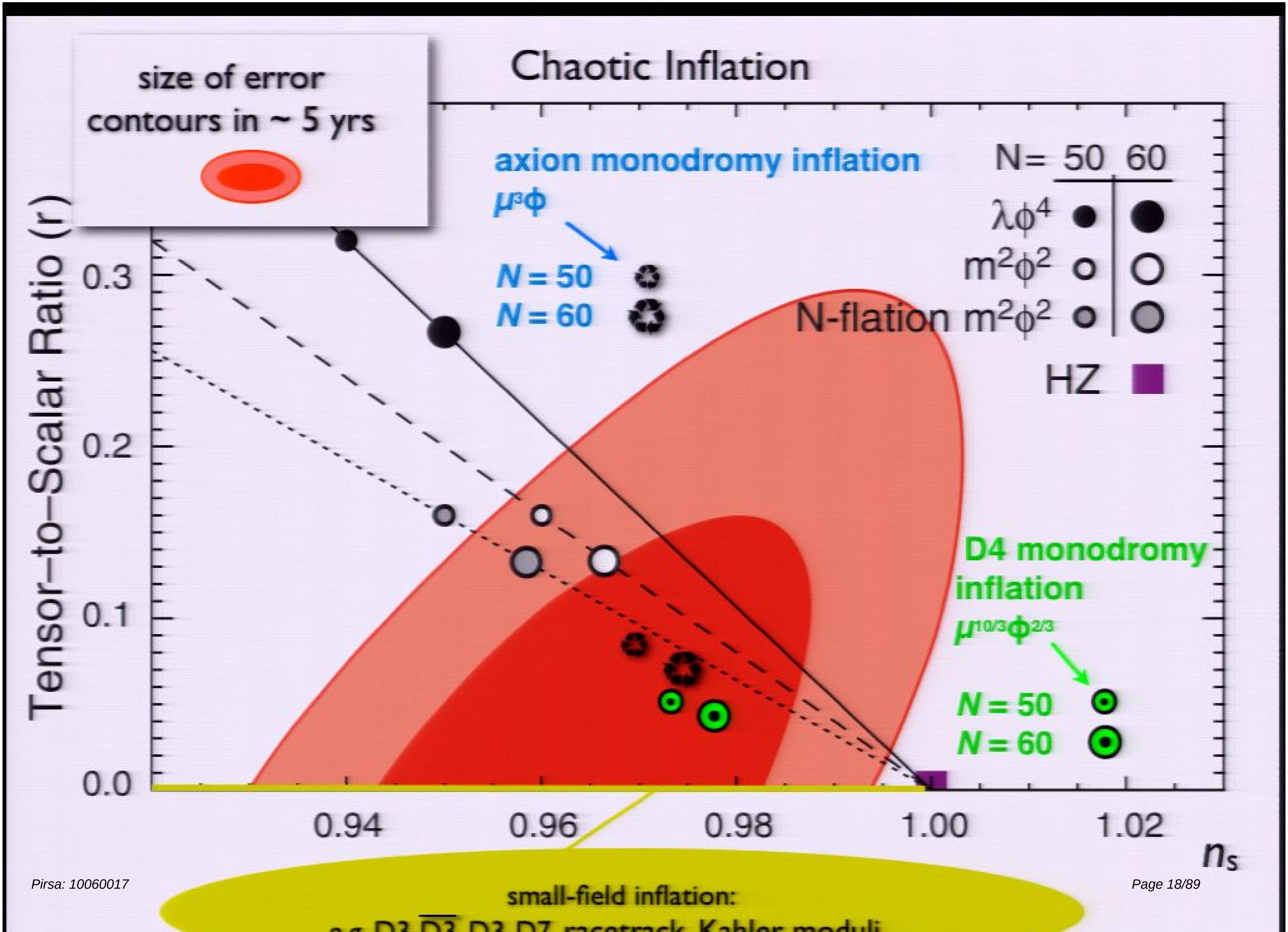
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the Kallosh-Linde problem ...

we are in 4D - string compactification ...

- we wish for low-energy supersymmetry - need to compactify internal 6 dimension on a Calabi-Yau manifold

- $\Rightarrow$ moduli: massless scalar fields, determining

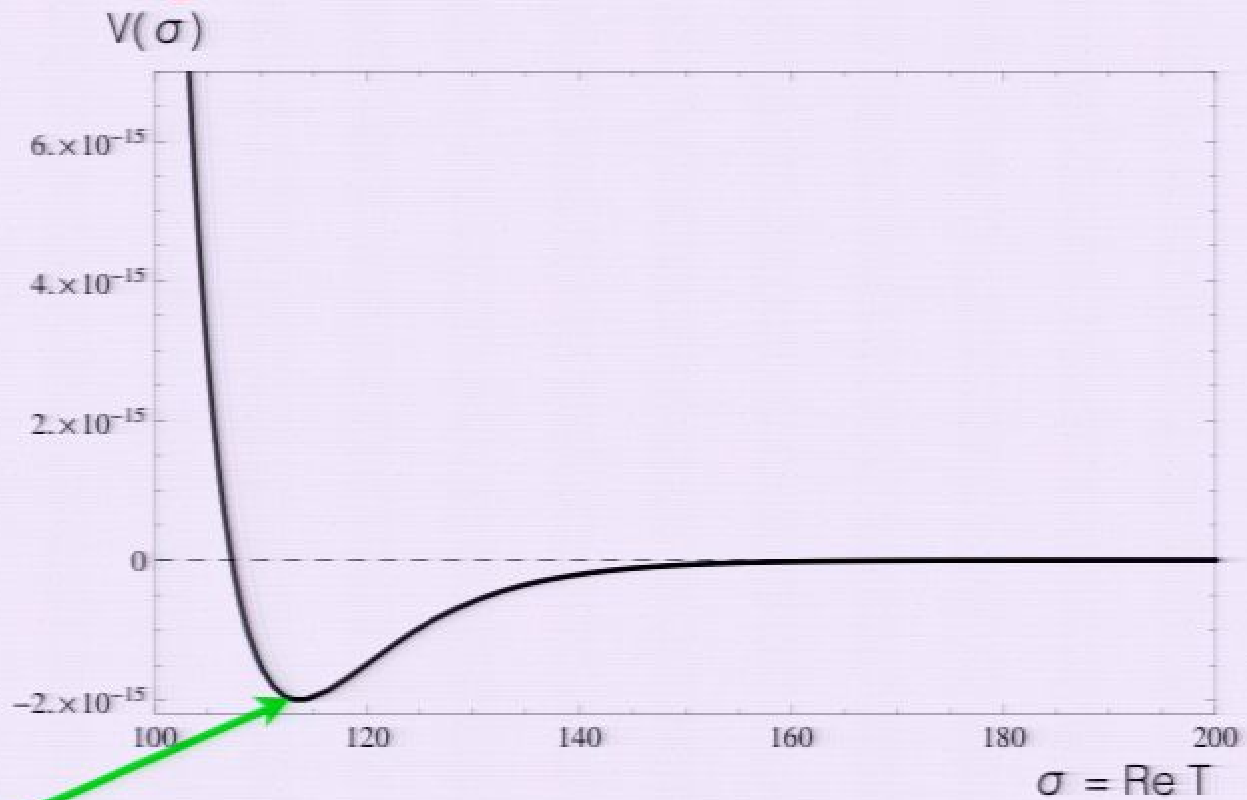
  size and shape of the CY

- $\Rightarrow$ one path to controlled compactification

  (KKLT) in IIB string theory:
  - fix the shapes with fluxes
  - fix the sizes with 1 instanton per size modulus

- single volume modulus case: an instanton balances against the non-*T* sector W_0 (e.g. from fluxes)

$$V(T) = e^K (K^{T\bar{T}} |D_T W|^2 - 3|W|^2)$$



$$K = -3\ln(T + \bar{T})$$

$$W = W_0 + Ae^{-aT}$$

fixes shape moduli

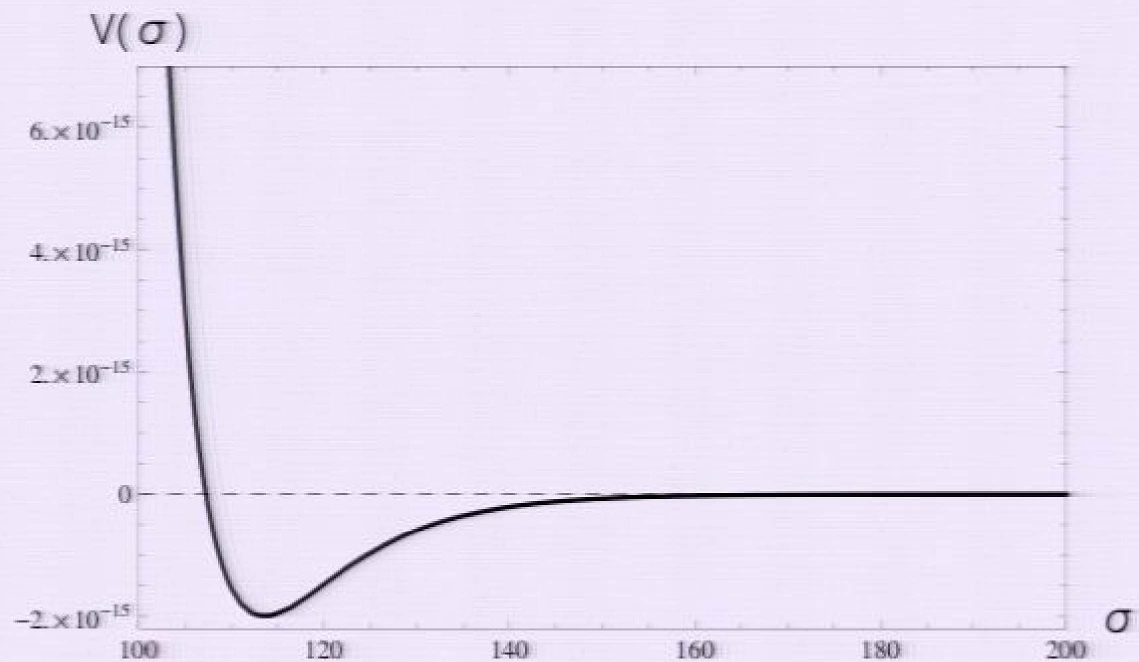
$$T_0 : \; D_T W(\varphi)|_{T_0} = 0$$

- inflationary sector generates a large positive vacuum energy
- by locality in the extra dimensions all energy forms can at most grow as fast as the volume
- Weyl rescaling into 4D Einstein frame - all energy forms scale as σ^{-3} = volume^{-2}
- $\Rightarrow$ all potential vanish at infinite volume &

  all positive energy states are metastable to *de-compactification*

- Einstein frame rescaling - SUSY breaking scales as inverse power of the volume σ = Re *T*

$$|V_{AdS}| = 3e^K |\langle W \rangle_0|^2$$



- Einstein frame rescaling - SUSY breaking scales as inverse power of the volume σ = Re *T*

$$V(\Phi) \sim e^K G^{\Phi\bar{\Phi}} |D_\Phi W|^2 \sim \frac{1}{\sigma^r}$$

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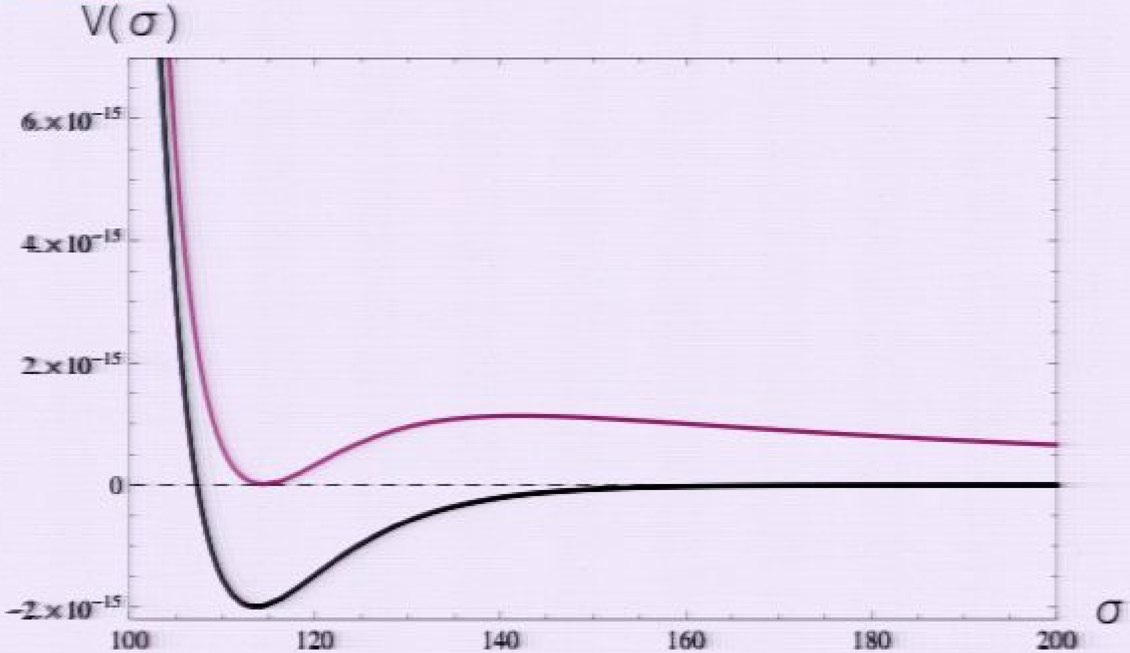
$$V_B \simeq |V_{AdS}|$$



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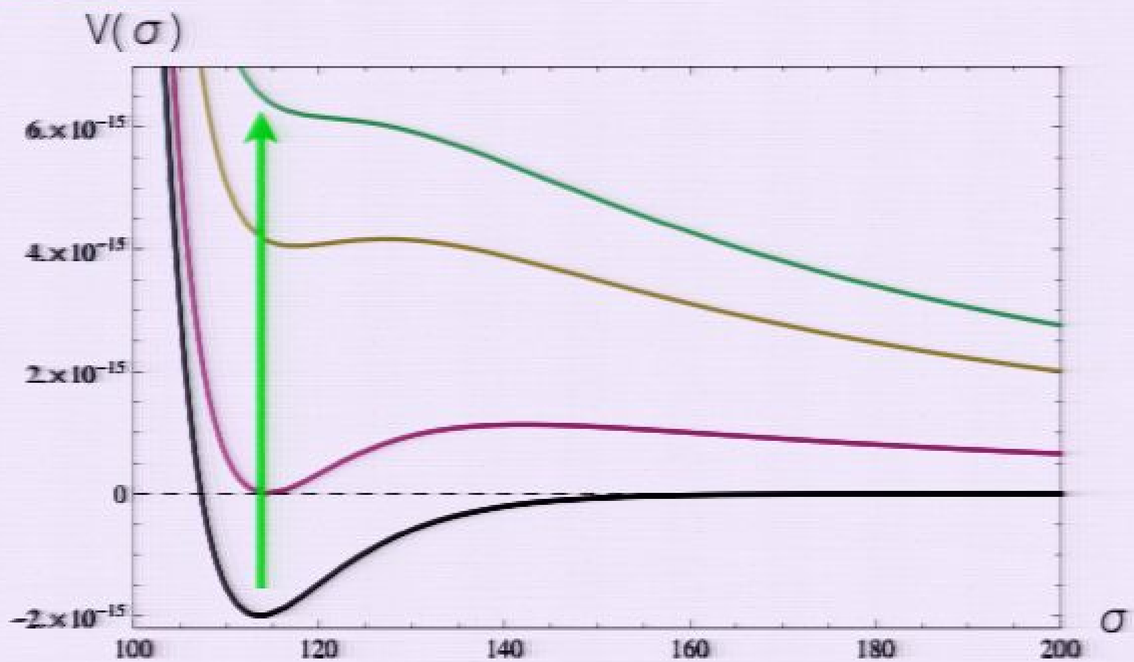


$$H^2 \lesssim \mathcal{O}(10) V_B \simeq \mathcal{O}(10) |V_{AdS}| \sim e^K |\langle W \rangle_0|^2 \sim m_{3/2}^2$$

overcoming KL ...

What to do ?

- decouple the barrier height from the (post-) inflationary uplifting: racetrack model of Kallosh & Linde, heavily fine-tuned at $O(m_{GUT}/m_W) \sim 10^{-13}$

- alternative: have the barrier height adjusting with the rolling inflaton!

- $\Rightarrow$ in f we have to adjust W_0 to adjust the barrier height

in more detail ...

- the depth of the AdS vacuum given by W_0 determines the barrier height induced by the post-/inflationary vacuum energy density

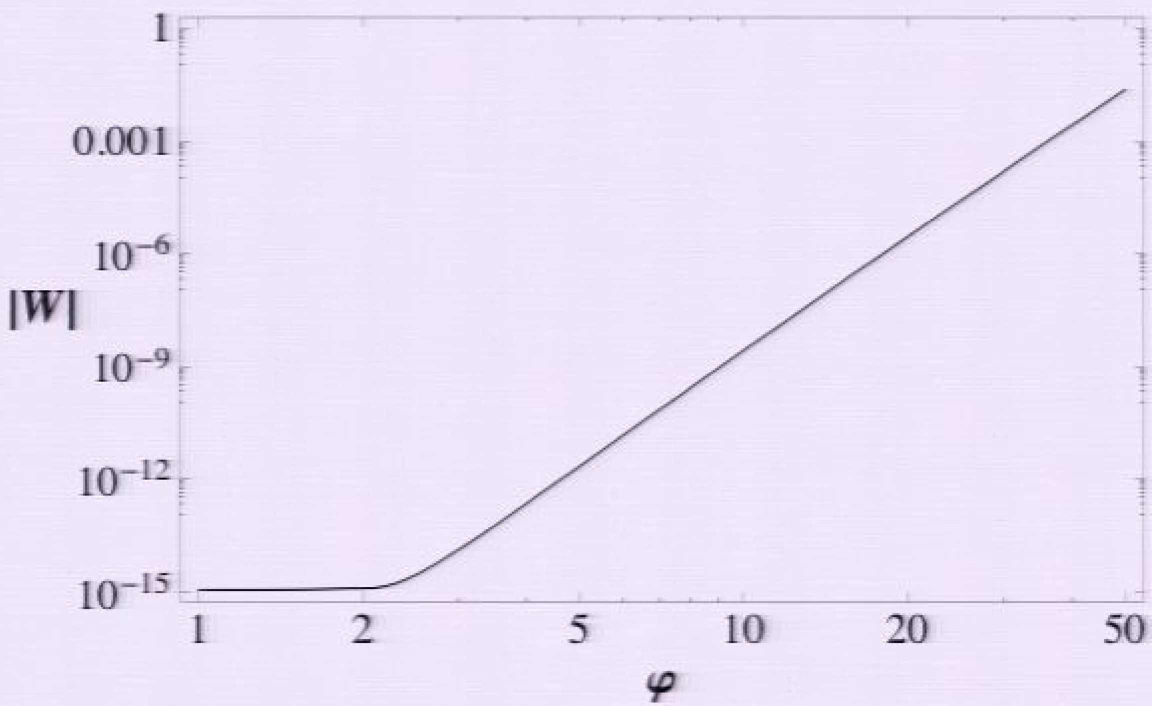
- for low-energy SUSY this leads to a very low barrier height, completely overrun by high-scale inflation

$\Rightarrow$ so if the flux-induced parameter W_0

controls the scales of the problem ...

- Who says, we cannot have W_0 being an adiabatic function of the inflaton?

$$W = W_{0,eff.}(\Phi) \qquad + Ae^{-aT}$$



- Let's try find simple models doing that ... However, in supergravity we cannot just rely on the inflation alone:

$$\frac{|F_\Phi|}{\sqrt{3}e^{K/2}|\langle W\rangle|} \approx \frac{n\alpha b\Phi^{n-1}}{\sqrt{3}(\alpha b\Phi^n + W_0)} \sim \frac{1}{\Phi}$$

- for a polynomial superpotential suitable for large-field inflation the potential slopes downward and goes *negative* ... So we probably have to get inflation from $F_X = F_X(\phi)$ from a 2nd field X

[Kawasaki, Yamaguchi & Yanagida '00]

- a simple setup which adjusts the barrier dynamically

$$K = \frac{1}{2}(\Phi+\bar{\Phi})^2 + X\bar{X} - \gamma(X\bar{X})^2 - 3\log(T+\bar{T})$$

$$W = W_0\, g(X) + \alpha\, f(X)\,\Phi^n + e^{-aT}$$

$$\text{with}:\ g(X) = 1 + \mathcal{O}(X) \quad \text{and} \quad f(X) = b + X + \mathcal{O}(X^2)$$

- this is t'Hooft natural, given that ϕ has R-charge 2/n and a shift symmetry in the Kähler potential:

$$\Phi = \eta + i\varphi \quad , \quad \varphi \to \varphi + C$$

- why do we need the 1st few terms in f and g, which are otherwise arbitrary?

- constant term in g gives us back the known KKLT-like post-inflation vacuum

- constant term in f we need to have W scaling adiabatically with ϕ

- the linear term in f in X we need to get that $F_X \sim W$, so that the potential slopes upwards ...

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- in this 4D *N*=1 supergravity the scalar potential reads

$$V = e^K (K^{\Phi\bar{\Phi}} |D_\Phi W|^2 + K^{X\bar{X}} |D_X W|^2 + K^{T\bar{T}} |D_T W|^2 - 3|W|^2) + \frac{C}{\sigma^2}$$

F_Φ F_X F_T

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F_Φ $\quad$ F_X $\quad$ F_T

- the last term can again be, e.g., a warped anti-D3 brane, lifting the *post*-inflationary vacuum to zero

- in the regime $\phi >> M_P$ and $X < M_P$ there is an attractor behaviour satisfying

$$F_X \sim \langle W \rangle \sim \alpha \Phi^n \quad , \quad F_\Phi \sim \frac{F_X}{\Phi} \quad , \quad F_T \sim \frac{F_X}{T}$$

- gives the inflaton potential to be

$$V_{inf.}(\varphi) \sim |F_X|^2 \sim \alpha^2 \varphi^{2n}$$

- and produces a mass term for X via

$$K^{X\bar{X}} = (1 - 4\gamma X\bar{X})^{-1} \simeq 1 + 4\gamma X\bar{X} \qquad \Rightarrow X \lesssim M_P$$

- Thus, we get a generalized KL-like constraint for the adiabatically adjusting VEV of W

$$\frac{\sqrt{|F_\Phi^2| + |F_X^2|}}{\sqrt{3}e^{K/2}|\langle W\rangle|} \sim \mathcal{O}(1)$$

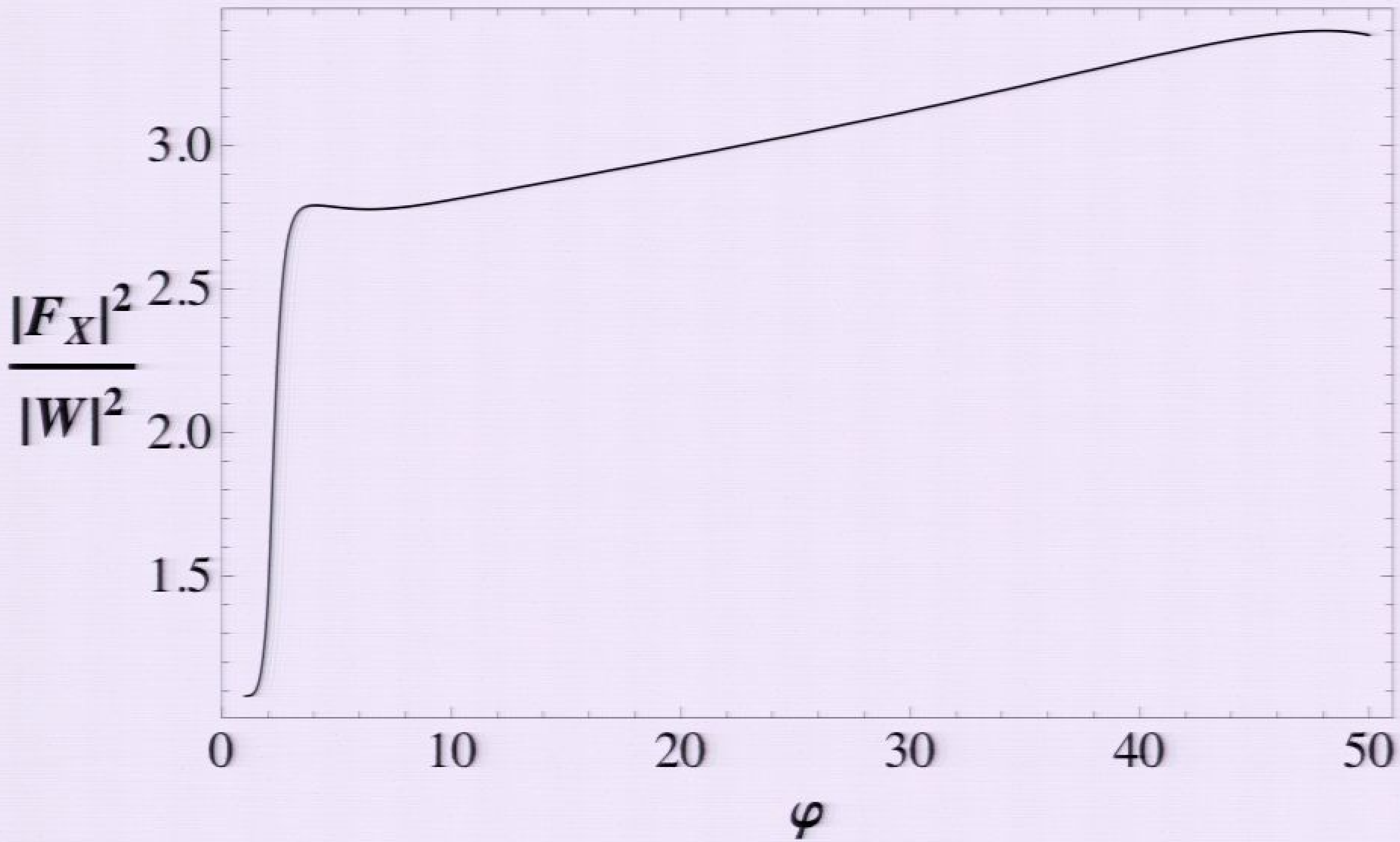
- at large $\phi >> M_P$ we get a scaling behavior

$$|F_X| \sim |W_{0,eff.}(\Phi)| \equiv |W_0 + \alpha(b+X)\Phi^n|$$

- which adjusts both (!) barrier height (controlled by W_0) and uplifting (controlled by F_X) dynamically such, that the minimum for T is never lost - if we adjust their ratio such that

$$|F_X|^2 \lesssim \mathcal{O}(10) 3 e^K |\langle W \rangle|^2$$

- this is done by choosing b



3.0
2.5
2.0
1.5
$\frac{|F_X|^2}{|W|^2}$
0
10
20
30
40
50
φ

- if this were all, the KL problem was fixed for good ...

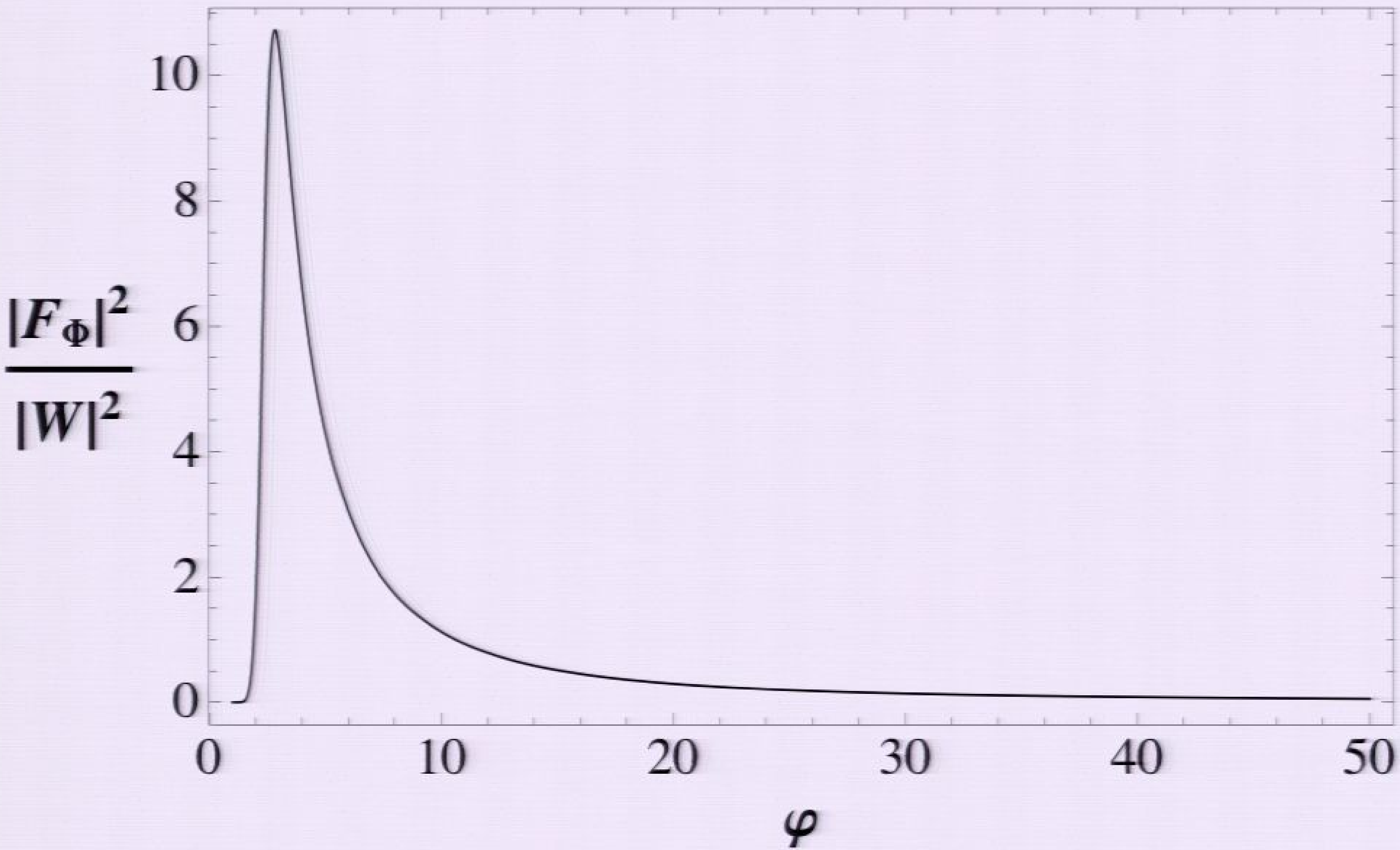
  However, at very small field ϕ such that $W_0 < \phi << M_P$, we have a regime $F_\phi > F_X$ with

$$\frac{|F_\Phi|}{|\langle W\rangle|} \approx \frac{n\alpha b\varphi^{n-1}}{\alpha b\varphi^n + W_0} \sim \frac{1}{\varphi}$$

  diverging, if W_0 vanishes

- at finite W_0 this ratio attains a maximum

$$\max\left(\frac{|F_\Phi|}{|\langle W\rangle|}\right) = (n-1)\left(\frac{\alpha b}{W_0(n-1)}\right)^{1/n}$$


$\frac{|F_\Phi|^2}{|W|^2}$
0
2
4
6
8
10
0
10
20
30
40
50
φ

- again, to guarantee stability of T, we must have

$$\max\left(\frac{|F_\Phi|}{|\langle W\rangle|}\right) \lesssim \mathcal{O}(10)$$

- for a given hierarchy W_0 / $W_{0,\text{eff.}}(\phi)$ and magnitude of density fluctuations δ this gives a *lower* bound on the power $2n$ of $V(\phi)$

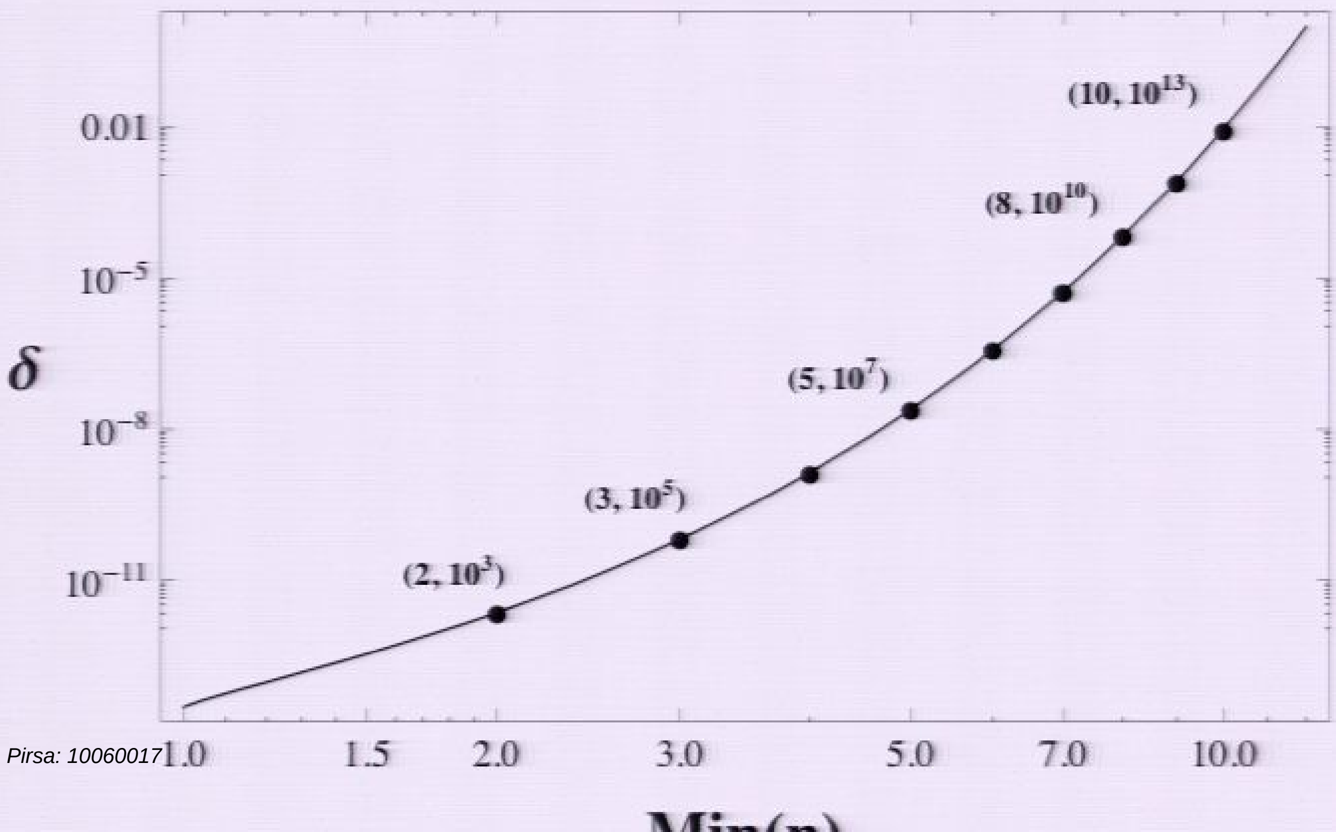
$$\alpha = \frac{10\sqrt{3}\pi n\delta}{\varphi_{60}^{n+1}} \quad , \quad \varphi_{60} = 2\sqrt{60(n-1)}$$

inflaton at 60 e-folds before inflation ends

$$(n-1)\left(\frac{10\sqrt{3}\pi b n\delta}{\varphi_{60}^{n+1}W_0(n-1)}\right)^{1/n} \leq 2\sqrt{3}$$

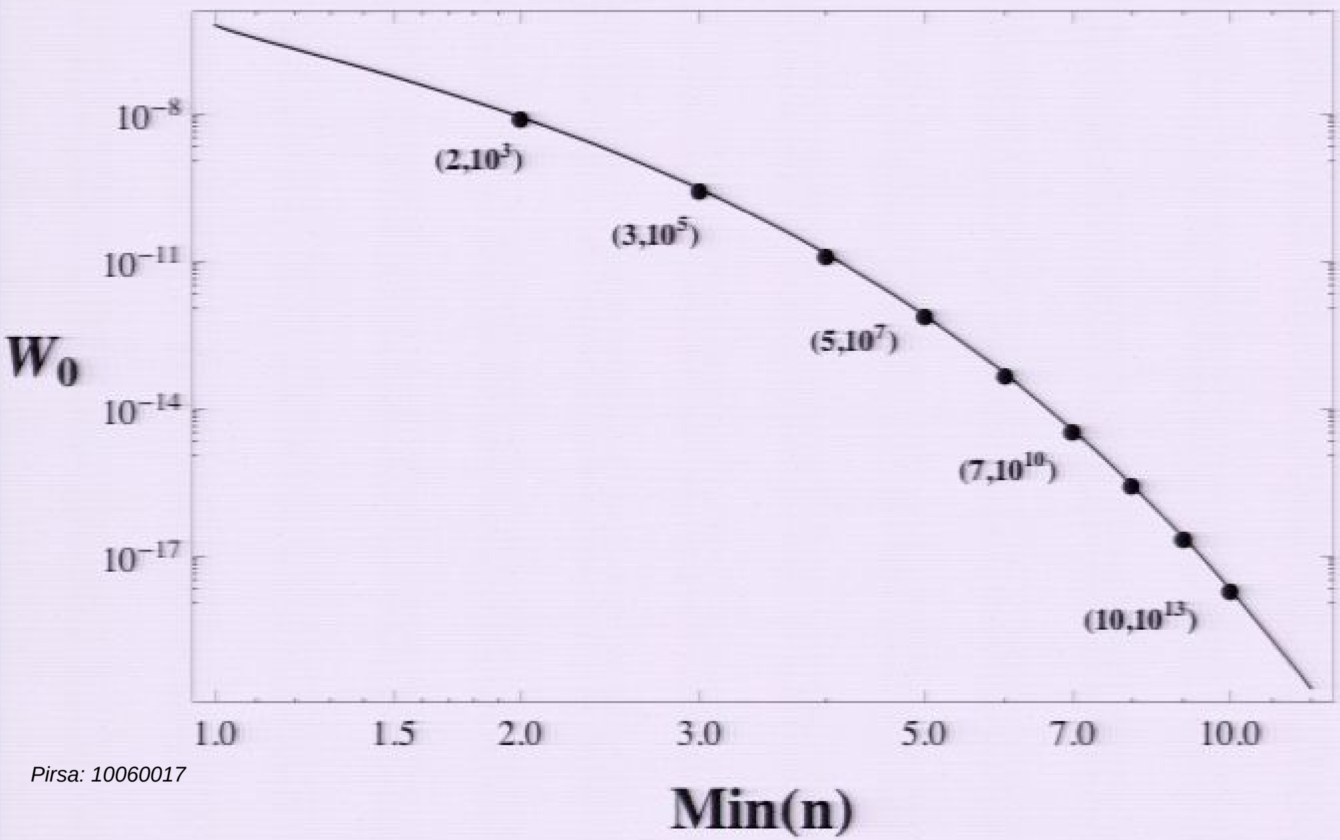
- thus we *horse trade*: if we wish to attain a given W_0 (TeV...) at the end of inflation, we can exchange n for δ

$$\delta \leq \frac{\left(\frac{2\sqrt{3}}{n-1}\right)^n \varphi_{60}^{n+1} W_0 (n-1)}{10\sqrt{3}\pi b n}$$



- *or*: if we wish to attain a given δ (2×10^{-5}) at ϕ_{60}, we can trade n for W_0 - and thus for the SUSY breaking scale after inflation

$$W_0 \geq \frac{10\sqrt{3}\pi n b \delta}{\left(\frac{2\sqrt{3}}{n-1}\right)^n \varphi_{60}^{n+1}(n-1)}$$

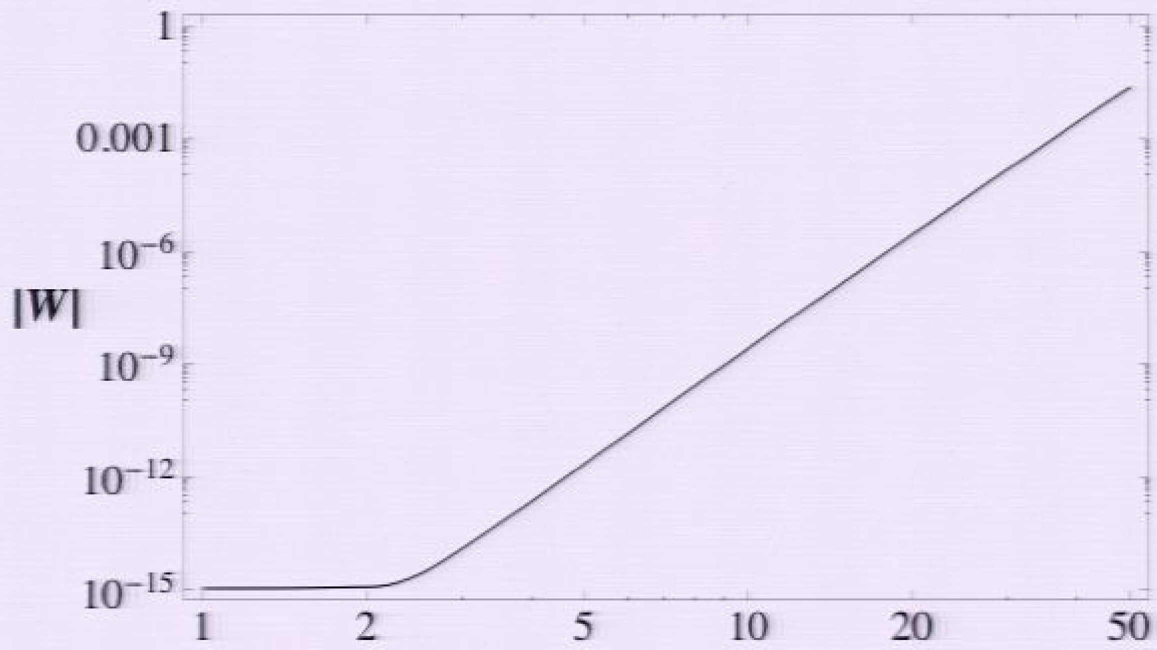


- a numerical example: evolution of all fields can be tracked, and used to prove the existence of adiabatic minima for X and T at all times - gives a huge hierarchy in W_0 during inflation ...

$$A = 1,\ a = \frac{2\pi}{10},\ W_0 = -10^{-15},$$

$$\alpha = 5 \times 10^{-19},\ b = \sqrt{2/5},\ n = 10,$$

and $\gamma = 2$



open questions ...

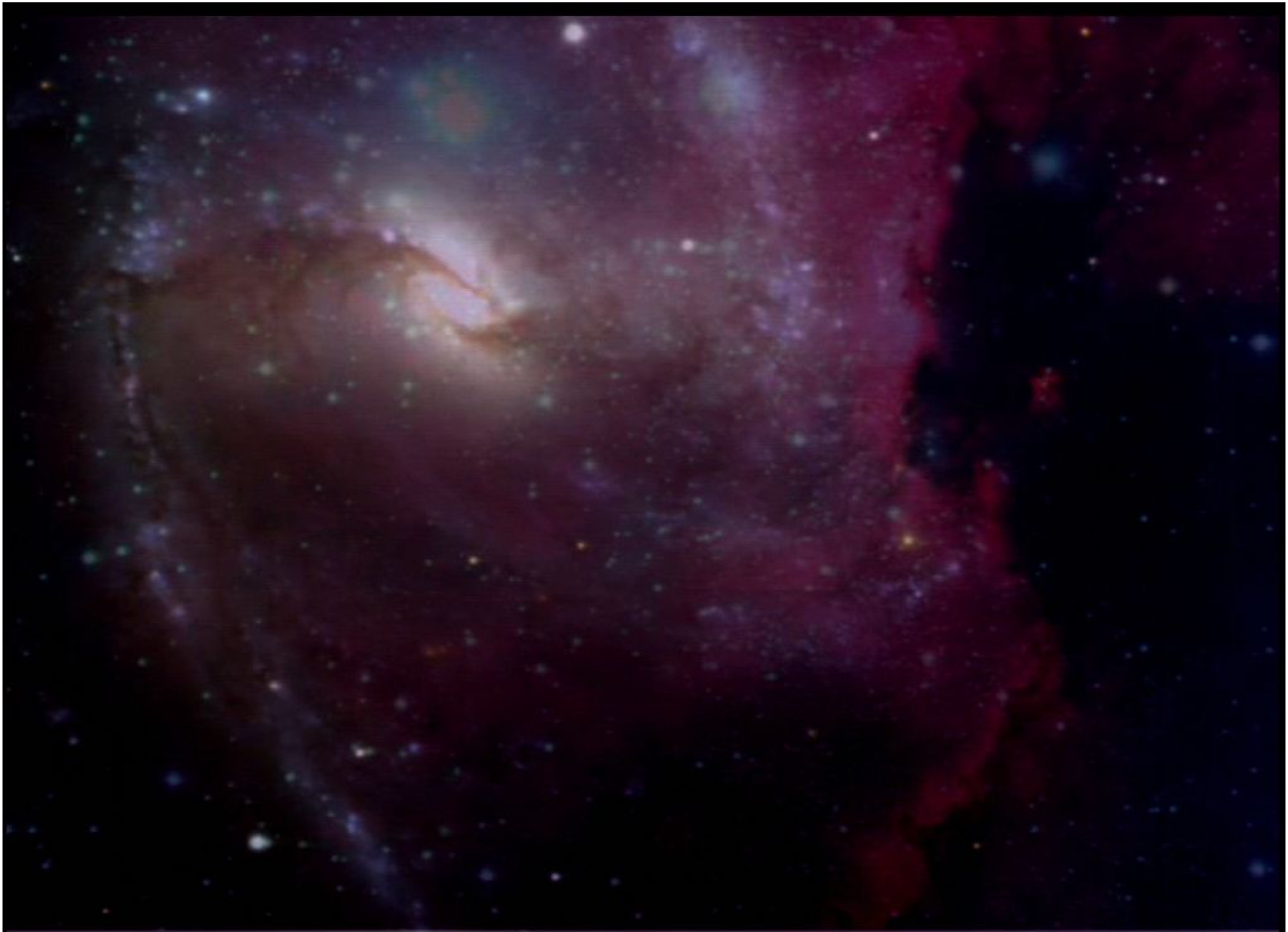
- how to get large-field inflation with large-ish powers in string theory?
  ... we know of (axion) monodromy inflation, which gives so far at most *linear* potentials ... [McAllister, Silverstein & AW '08/'09]

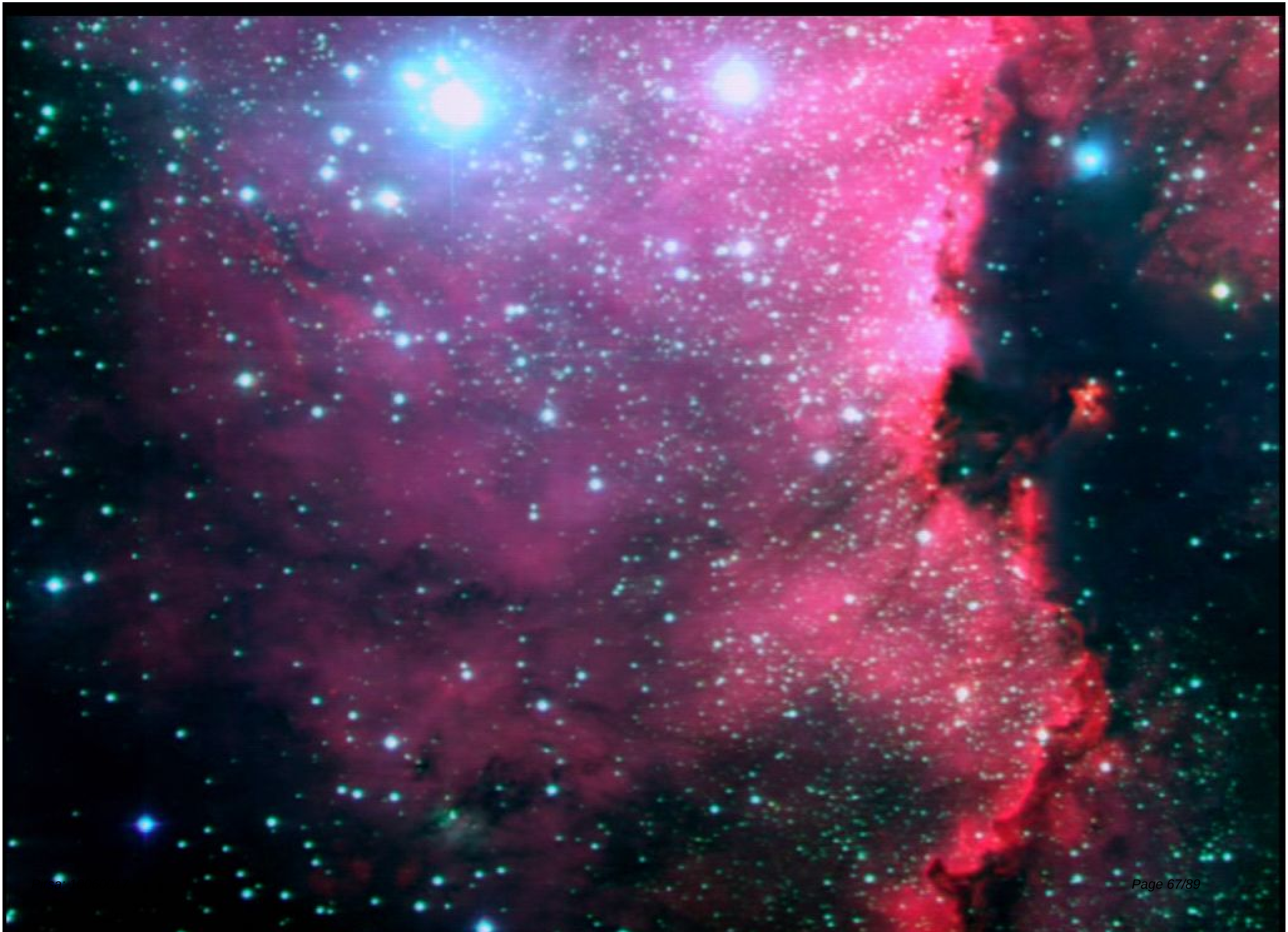
- the horse trading could be presumably loosened by modifying the exit in a hybrid-like fashion ...

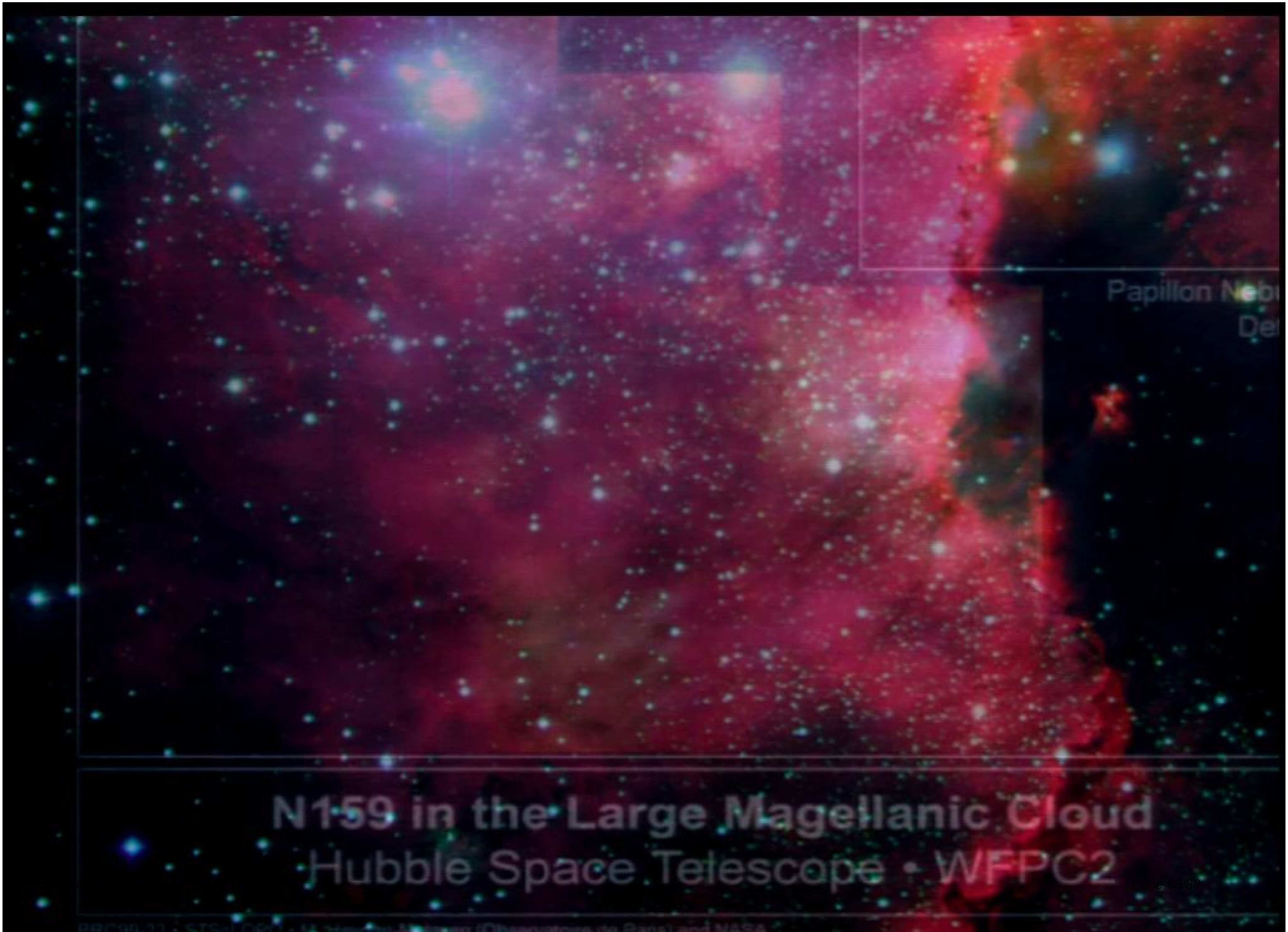
- small field models using the same basic mechanism?

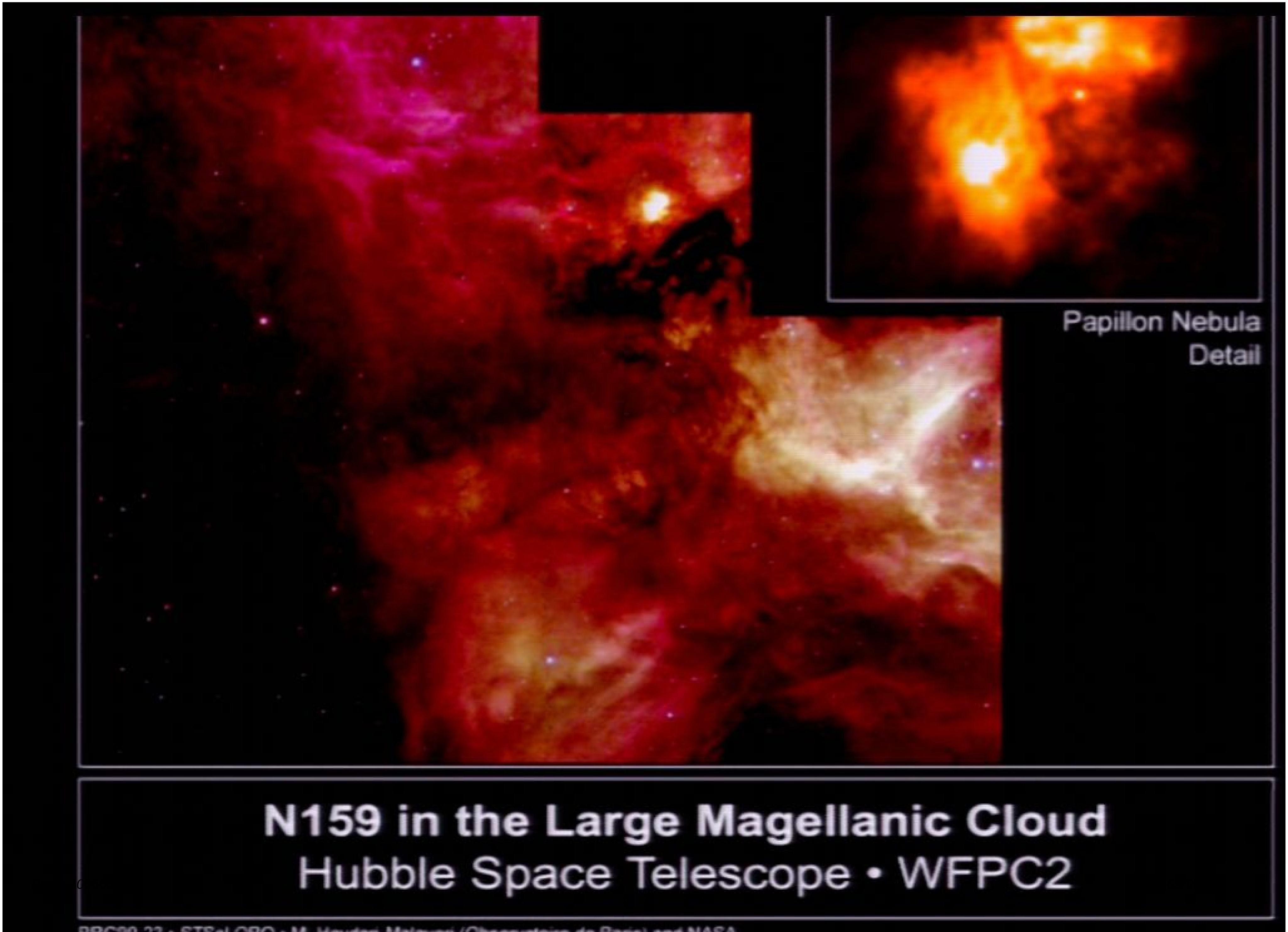
[redacted]

IR Colour Composite of RCW38 Region






N159 in the Large Magellanic Cloud
Hubble Space Telescope • WFPC2



Papillon Nebula
Detail
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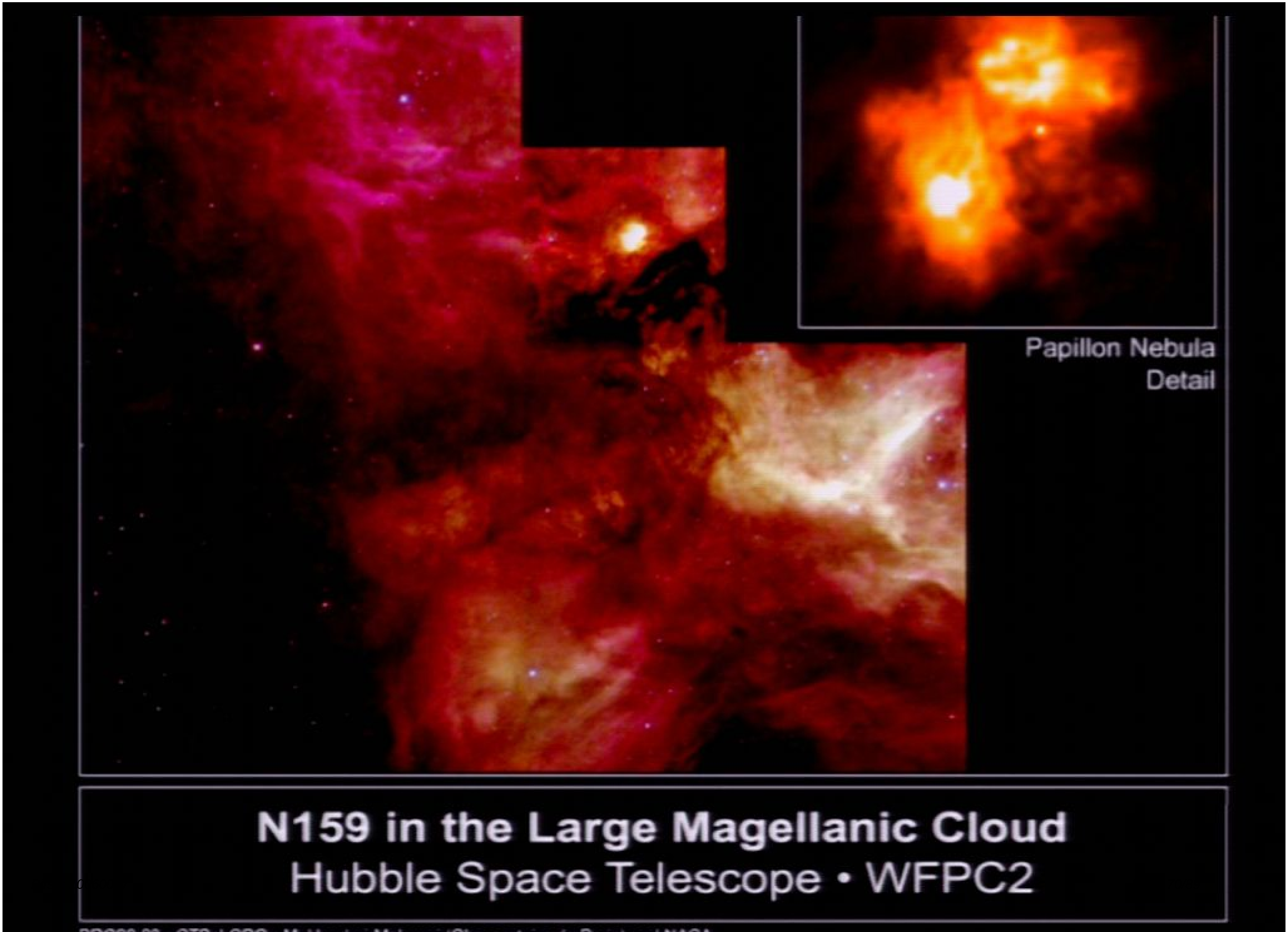
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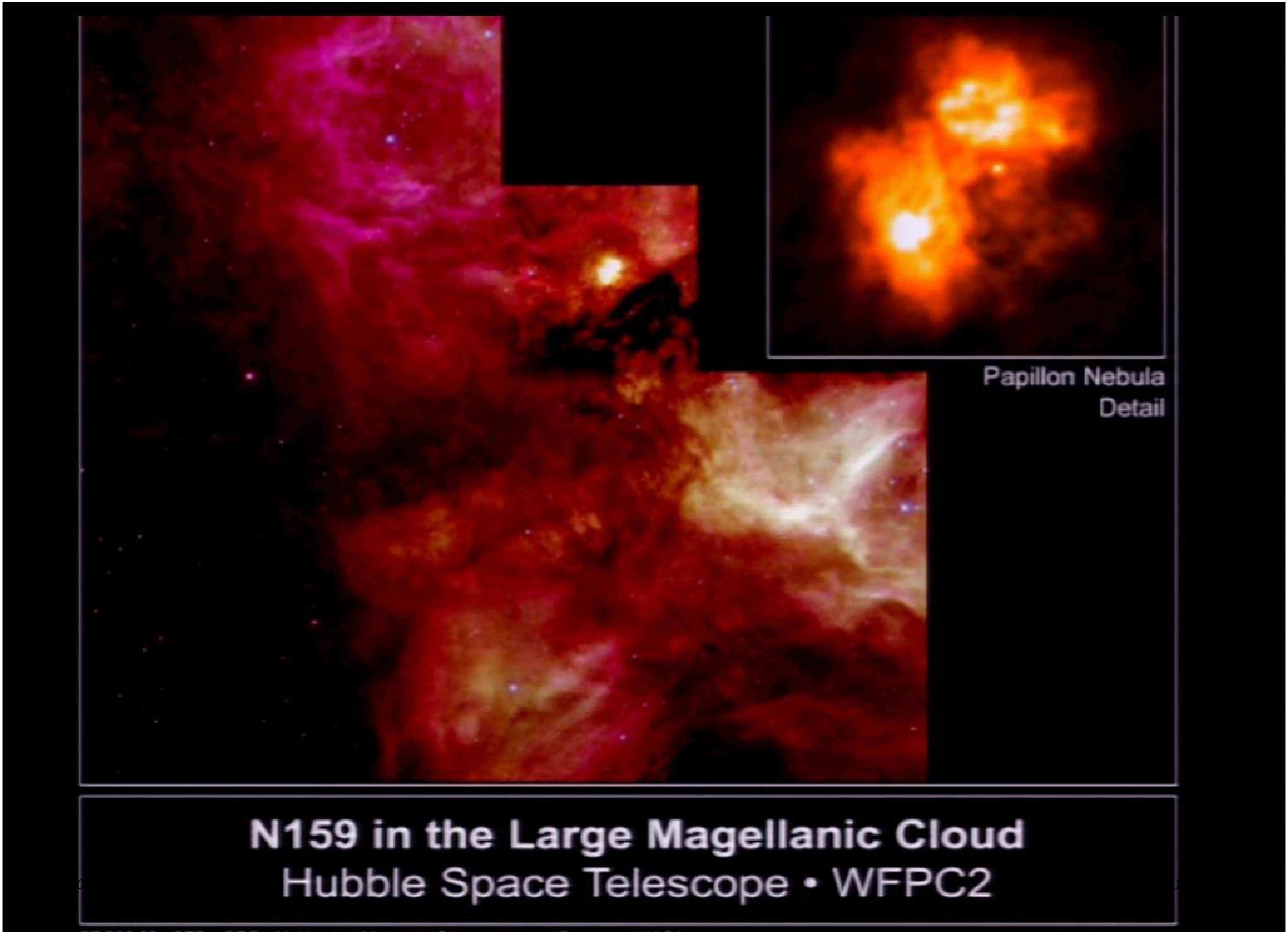
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