A New Screening Mechanism and its Cosmological consequences

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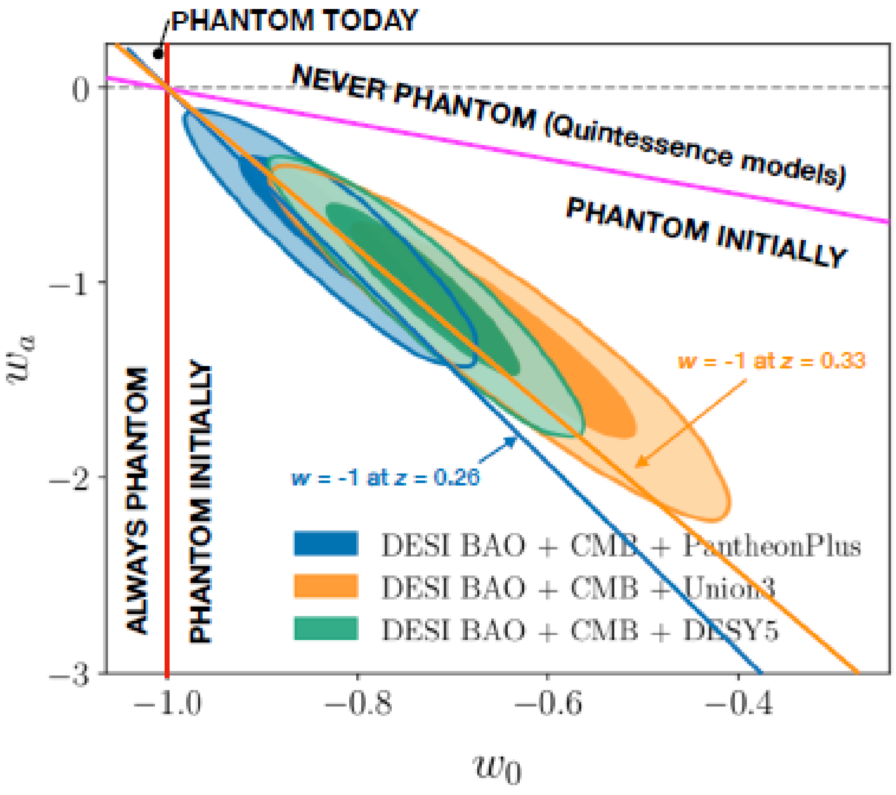
Collaboration with C. van de Bruck(Sheffield), C. Burgess (Perimeter), A. Davis (Cambridge), F. Quevedo (Cambridge) and A. Smith (Sheffield).



E pur si muove !

$$\underbrace{w_0 = -0.827 \pm 0.063 \qquad w_a = -0.75^{+0.29}_{-0.25}}_{\text{DESI + CMB + Pantheon+} \implies 2.5\sigma}$$

$$\underbrace{w_0 = -0.64 \pm 0.11 \qquad w_a = -1.27^{+0.40}_{-0.34}}_{\text{DESI + CMB + Union3} \implies 3.5\sigma}$$

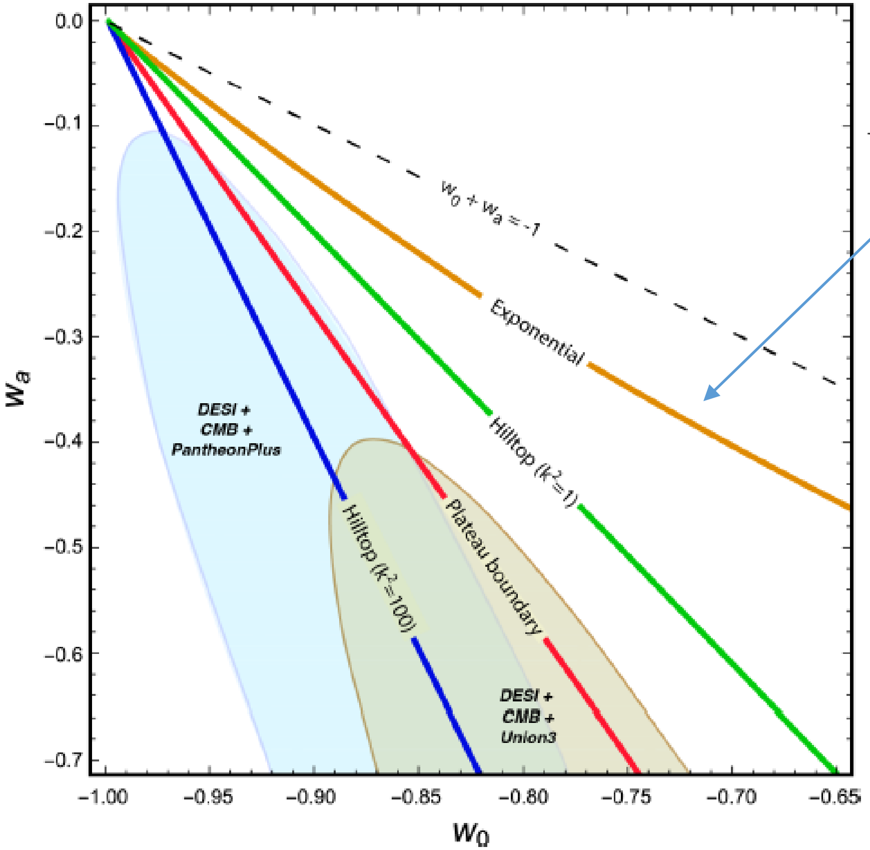
$$\underbrace{w_0 = -0.727 \pm 0.067 \qquad w_a = -1.05^{+0.31}_{-0.27}}_{\text{DESI + CMB + DES-SN5YR} \implies 3.9\sigma}$$

w_a

$$\omega(z) = \omega_0 + \omega_a \frac{z}{1+z}$$

Simple models can reproduce the data!

Is it expected?

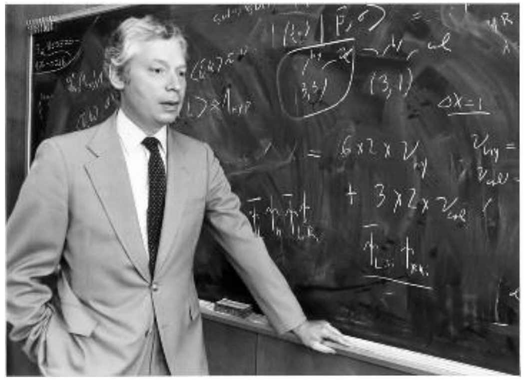


$V(\phi) = V_0 e^{-\lambda\phi/m_{\rm Pl}}$

The dark energy scale is in the ***pico-eV range***: apparent fine-tuning compared to standard model scales.

$$\xcancel{\delta\rho_\Lambda = M^4, \quad M \sim 100\mathrm{GeV}}$$

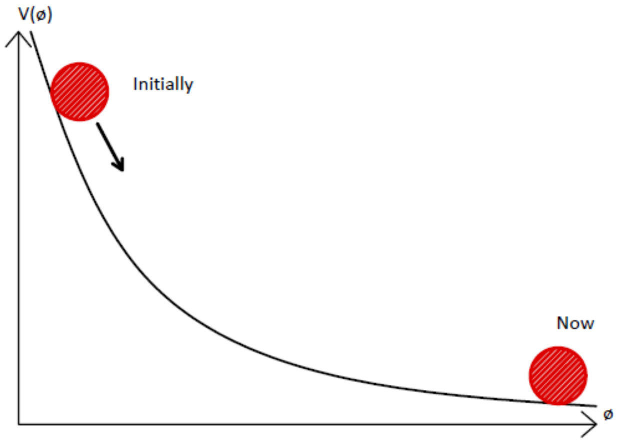
Weinberg's theorem states that there is no non-fined tuned vacuum in a 4d quantum field theory respecting **Poincare invariance**.



Dynamical configurations

↓

Dark energy



Scalar field rolling down its potential

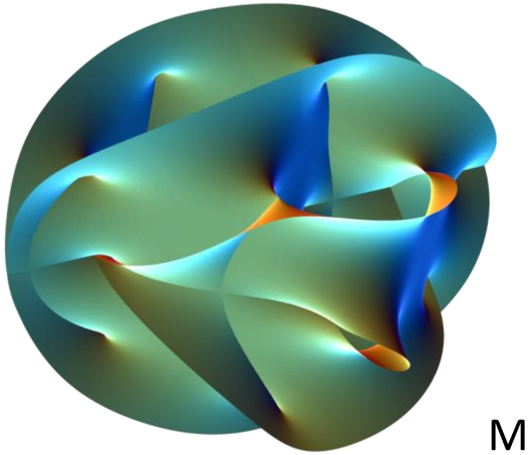
Lorentz invariance implies the existence of a dark energy field which can be seen as the Goldstone model for the breaking of time translation invariance.

The most important stringy ***conjectures*** for dark energy are:

- ✓ ***The de Sitter conjecture***: a pure vacuum energy with no dynamics is not compatible with string theory.

In an ideal world, string theory or any other version of quantum gravity would be finite so the vacuum energy could be calculable. Not the case in ordinary Quantum Field Theory.

- ✓ ***The vacuum conjecture:*** Empty space-time is described by the dynamics of at least one scalar field with a potential such that



$$|\frac{dV}{d\phi}| \geq c\frac{V}{m_{\rm Pl}}$$

$$c = \mathcal{O}(\sqrt{2})$$

This forbids very flat potentials. This favours runaway potentials where the field is a "moduli".

Moduli could be "sizes" of extra-dimensions

Some expected features:

- Dark energy is determined by the position of the field now:

$$3\Omega_\Lambda H_0^2 m_{\rm Pl}^2 = V(\phi_{\rm now})$$

- The field is ***extremely light***:

$$m_\phi^2 = \frac{d^2V}{d\phi^2}|_{\rm now} \sim \frac{V_{\rm now}}{m_{\rm Pl}^2} = 3\Omega_\Lambda H_0^2$$

Mass of the order
of the Hubble rate

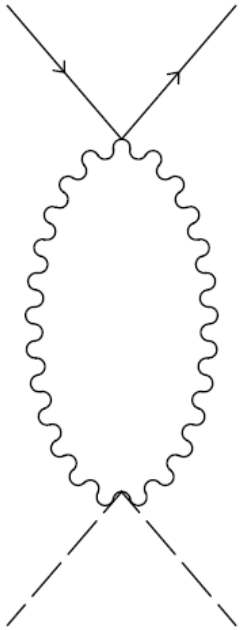
$$H_0 \sim 10^{-42}\ \text{GeV}$$

Problem: The coupling to matter

small

$$\beta \sim \frac{H^2}{m_{\rm Pl}^2}\int \frac{d^4p}{p^4}$$

Divergent



$$\mathcal{L} \supset -\frac{\beta}{m_{\rm Pl}} m_\psi \phi \bar{\psi}\psi$$

Yukawa interaction similar to the Higgs interaction to matter.

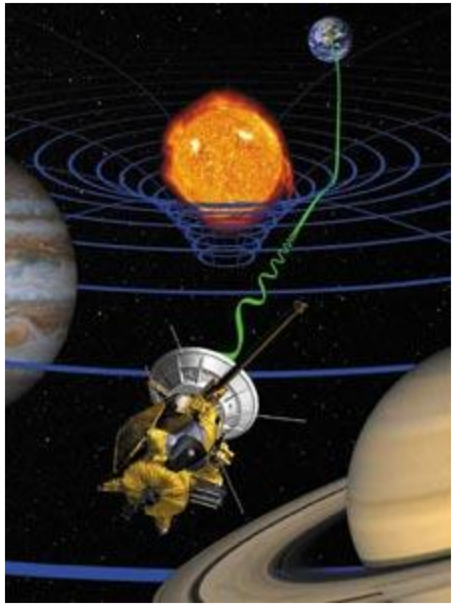
No ***reason*** to assume $\beta = 0$

Deviations from Newton's law are parametrised by:

$$\Phi = -\frac{G_N M}{r}(1 + 2\beta^2 e^{-r/\lambda})$$

For long range forces with large λ, the tightest constraint on the coupling β comes from the Cassini probe measuring the Shapiro effect (time delay around a big object: the Sun):

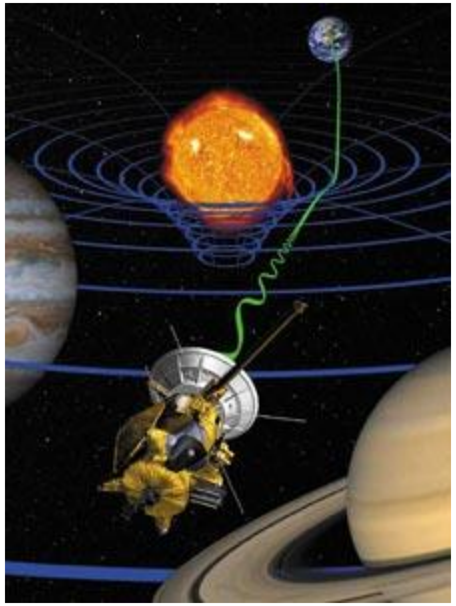
$$\beta^2 \leq 4 \cdot 10^{-5}$$



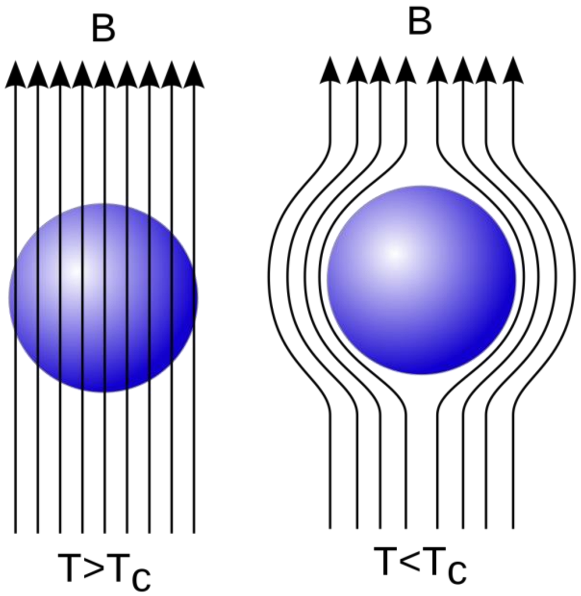
Bertotti et al. (2004)

Two options:

Symmetry: the shift symmetry of a Goldstone boson prevents such a coupling BUT the symmetry is broken by the potential so the problem is reintroduced!



Screening



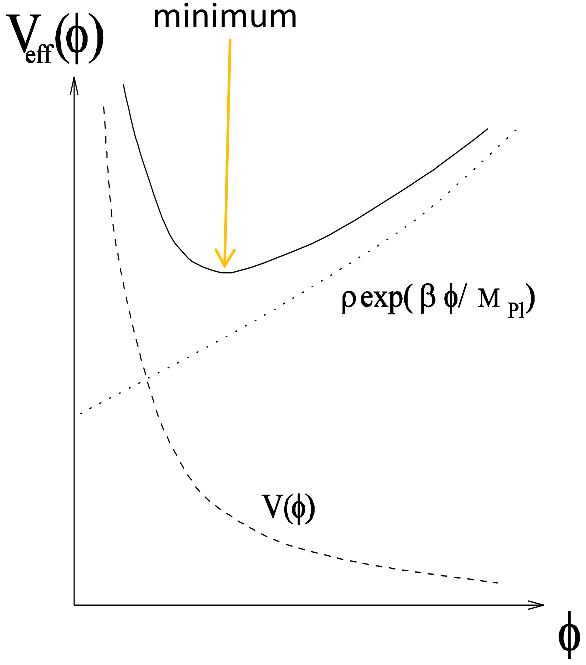
Analogous to the Meissner effect in superconductors: the inside of the solar system is free of scalar field lines. → No gradient=no fifth force

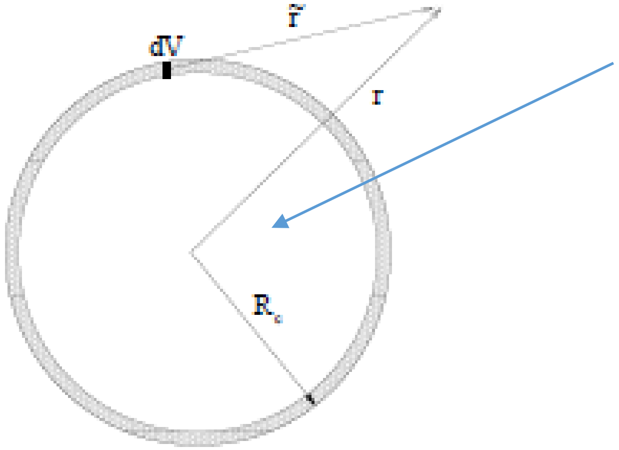
Chameleon Screening

When coupled to matter, scalar fields have a ***matter dependent effective potential***

$$\boxed{V_{eff}(\phi) = V(\phi) + \rho_m(A(\phi) - 1)}$$

Environment dependent minimum





Large mass inside object

$$m_{\rm in} R \gg 1$$

The field generated from deep inside is Yukawa suppressed. Only a *thin shell* radiates outside the body. Hence suppressed scalar contribution to the fifth force.

Vainshtein Mechanism in a nutshell

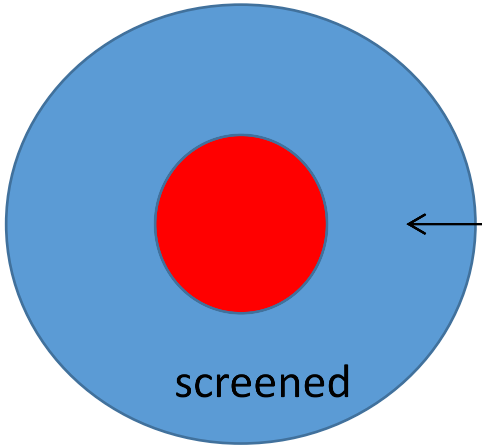
We can use a simple example with higher derivatives:

$$\mathcal{L}_\phi = -\frac{1}{2}(\partial\phi)^2 - \frac{1}{2\Lambda^3}(\partial\phi)^2\Box\phi + \frac{\beta\phi}{M_P}T\,.$$

Non-linearity

$$\Lambda^3 = H_0^2 m_{\mathrm{Pl}}$$

Very low cutoff scale!



$$\frac{F_\phi}{F_N} = 2\beta^2\left(\frac{r}{R_V}\right)^{3/2}\,.$$

$$R_V = \left(\frac{\beta M_c}{2\pi M_P}\right)^{1/3}\frac{1}{\Lambda}\,.$$

Well inside the Vainshtein radius, Newtonian gravity is restored. Well outside, gravity is modified.

The Vainshtein radius is very large for stars, typically 0.1 kpc for the sun.

These theories are beyond normal effective theories

$$V(\phi) = \Lambda_0^4 + \boxed{\frac{\Lambda^{n+4}}{\phi^n}}$$

Inverse power laws… and needs a cosmological constant

$$\mathcal{L}_\phi = -\frac{1}{2}(\partial\phi)^2 - \frac{1}{2\Lambda^3}(\partial\phi)^2\Box\phi + \frac{\beta\phi}{M_P}T \; .$$

Vainshtein when: $\boxed{\Box \geq H_0^2}$

Need to work beyond the validity of the derivative expansion…. (possible way out: non-renormalisation theorems)

Bringing modified gravity back to the fold:

Take heed from successful physics models:

The standard model of particle physics:

$$\mathcal{L} \supset \frac{\mu^2}{2}H^2 - \frac{\lambda}{4}H^4 + yH\bar{\psi}\psi$$

Simplest lowest order effective theory coupling Higgs and fermions

Inflation :

$$\mathcal{L} = \frac{R}{16\pi G_N} + cR^2 + \ldots$$

Starobinski: simple curvature expansion around GR…

Multi-field dark energy sector

$$\mathcal{L} = -\frac{1}{2} g_{ij}(\phi^k)\partial_\mu \phi^i \partial^\mu \phi^j - V(\phi^i)$$

Your favourite dark energy potential with no wacky potential in it.

Screening can only happen with more than one field and a non-trivial σ-model metric.

Typical example:

$$\phi$$

Dilaton of broken scale invariance

$$a$$

Pseudo-Goldstone axion of broken symmetry

Runaway dilaton model:

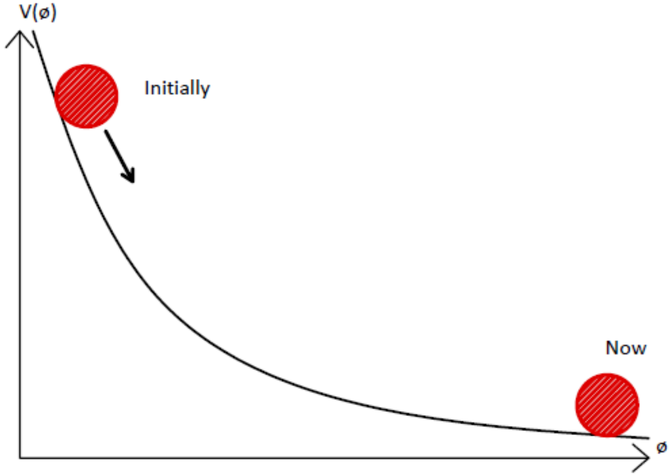
Focus on simplest potential with runaway behaviour:

Klein-Gordon equation

$$V(\phi) = V_0 e^{-\lambda\phi/m_{\mathrm{Pl}}}$$

$$\ddot{\phi} + 3H\dot{\phi} + \frac{\partial V}{\partial \phi} = 0$$

$$\gamma = \frac{2}{\lambda}, \quad \alpha = \frac{2}{\lambda^2}$$



Friedmann:

$$H^2 = \frac{\frac{1}{2}\dot{\phi}^2 + V(\phi) + \rho}{3m_{\mathrm{Pl}}^2}$$

matter

$$\omega_\phi = -1 + \frac{\lambda^2}{3}$$

Initially dark energy is subdominant and then starts dominating when matter becomes very small.

$$\phi = \phi_\star + \gamma m_{\mathrm{Pl}} \ln \frac{t}{t_\star}$$

$$a = a_\star(\frac{t}{t_\star})^\alpha$$

Yoga models:

Large so no acceleration without the U prefactor.

$$V(\phi) = U(\phi)e^{-\lambda\phi/m_{\mathrm{Pl}}}$$

Quadratic with a minimum.

The axio-dilaton system

$$\mathcal{L} \supset -\frac{1}{2}((\partial\phi)^2 + W^2(\phi)(\partial a)^2)$$

The axio-dilaton case corresponding to the volume modulus of compactifications from 10d to 4d:

$$K = -3m_{\rm Pl}^2 \ln(T + \bar{T})$$

The volume modulus can be decomposed in:

$$T = e^{\phi} + ia \Longrightarrow e^{K/6m_{\rm Pl}^2} = e^{-\phi/2}$$

axion

dilaton

$$W(\phi) = e^{-\phi}$$

Dilaton coupling to matter
of order one!

The axio-dilaton system

$$\mathcal{L} \supset -\frac{1}{2}((\partial\phi)^2 + W^2(\phi)(\partial a)^2)$$

Choose W to have a minimum:

$$W^2(\phi) = 1 + \frac{(\phi - \phi_\star)^2}{2\Lambda_\phi^2}$$

The axion is chosen to have a simple potential:

$$V(a) = \frac{1}{2} m_a^2 (a - a_+)^2 + \rho_m \frac{(a - a_-)^2}{2\Lambda_a^2}$$

QCD-inspired…

The Klein-Gordon equation for the axion is simply:

$$\nabla_\mu(W^2(\phi)\partial^\mu a) = \frac{1}{f^2}\partial_a V_{eff}$$

$$V_{\rm eff}(a) = V(a) + \rho(x)U(a)$$

Matter density

The QCD example:

$$V_{\rm QCD}(a) = -\Lambda^4_{\rm QCD}(1 - \frac{4m_u m_d}{(m_u+m_d)^2}\sin^2(\frac{a}{2}))^{1/2}$$

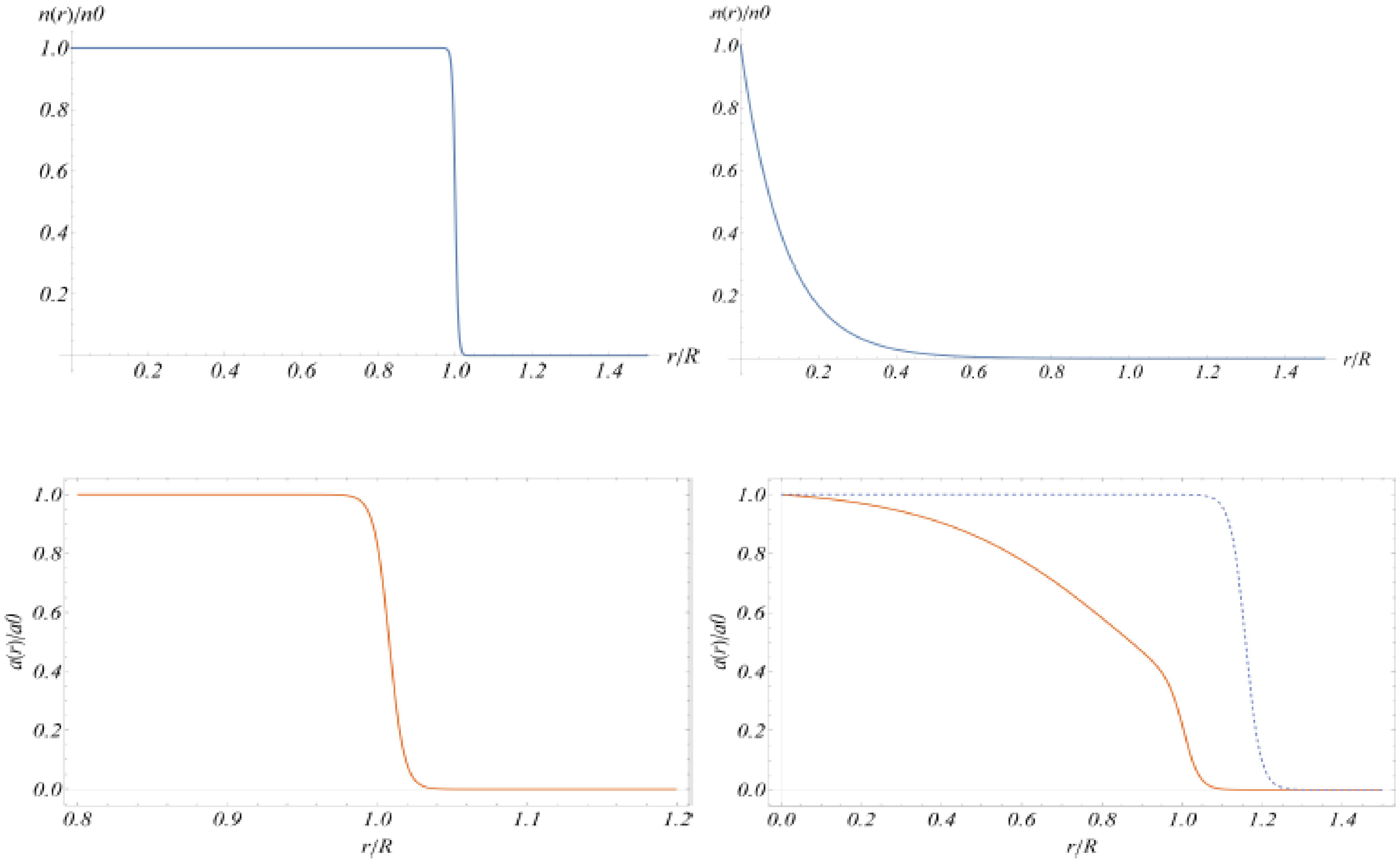
$$U(a) = \sigma_B \cos\frac{a}{2}$$

Minimised at a=0 in vacuum to solve the QCD CP problem.

Of order of the QCD scale

Minimised at a=π

The axion profile depends on the mass of the axion inside and outside matter.



$$m_{\text{in}} R \gg 1$$

Size of the object.

$$\mathcal{L} \supset -\frac{1}{2}((\partial\phi)^2 + W^2(\phi)(\partial a)^2)$$

The large gradients impose that the dilaton does not vary much locally:

$$\phi \approx \phi_\star$$

Screening !

Large gradients

Minimised when W is minimal

The Klein-Gordon for the dilaton is:

Coupling constant to matter.

$$\Box\phi = WW'(\partial a)^2 + \frac{\beta}{m_{\mathrm{Pl}}^2}\rho$$

Axion driven “potential”

Far away from the body we expect:

$$\phi = \phi_\infty - \frac{L m_{\mathrm{Pl}}}{r}$$

$$\boxed{L = 2\beta G_N M}$$

In the absence of screening.

Local value of the field in the environment

The delta function as a source term implies a jump of the derivative of the dilaton:

$$\phi'_{\rm out}(R) - \phi'_{\rm in}(R) = (\frac{WW'}{2\ell})_{r=R}(a_+ - a_-)^2$$

$$Lm_{\rm Pl} = R^2\phi'_{\rm out}(R)$$

Needs to be negative to reduce the scalar charge of the object

The scalar charge L is determined by the competition between the different energy sources for the scalar field profile. This competition will select the **local value of the field** and determine the scalar charge.

In the axio-dilaton case, the solar system is small enough that the local energy is dominated by the local dynamics which differ from the cosmological one.

$$E_{\rm kin} = 2\pi \int_0^\infty dr\; r^2 (W^2 (a')^2 + (\phi')^2)$$

Gives a surface contribution

$$E_{\rm kin,a} = \frac{\pi}{\ell} R^2 W^2 (r = R)(a_+^2 - a_-^2)$$

Dependent on the local value

$$E_{\rm kin} = E_0 + \frac{L^2}{4 G_N R}$$

$$L = 2\beta G_N M + R^2 (\frac{W W'}{2\ell})_{r=R} \frac{(a_+ - a_-)^2}{m_{\rm Pl}}$$

Having expanded W close to a minimum, the effective coupling to matter L can be obtained by minimising the energy with respect to the local value:

$$\beta_{\rm eff} \equiv \frac{L}{2G_N M} = \frac{\beta}{1+\frac{R}{\ell}\frac{(a_+-a_-)^2}{4\Lambda_\phi^2}}$$

Screening takes place as:

$$\frac{\ell}{R} \ll 1 \Rightarrow \frac{\beta_{\rm eff}}{\beta} = \mathcal{O}(\frac{\ell}{R}) \ll 1$$

The dilaton is screened thanks to the large mass of the axion.

Exponential Screening:

Screening also takes place when:

$$W^2(\phi) = e^{-\xi\phi/m_{\rm Pl}}$$

The coupling to matter becomes:

$$\beta_{\rm eff} = \beta \frac{R}{2\xi G_N M}$$

Screening in solar system:

$$\xi \gtrsim 10^9 \Rightarrow \frac{m_{\rm Pl}}{\xi} \lesssim 10^9 {\rm GeV}$$

Cosmology

Two Main Effects:

- **Early Dark Energy**

- **Influence on the growth of structure**

Early dark energy for free!

$$V(a) = \frac{1}{2} m_a^2 (a - a_+)^2 + \rho_m \frac{(a - a_-)^2}{2\Lambda_a^2}$$

$\rho_m \gg \Lambda_a^2 m_a^2$

$$\boxed{V(a_-) = \frac{1}{2} m_a^2 (a_+ - a_-)^2}$$

Early dark energy

A fraction of added matter

$\rho_m \ll \Lambda_a^2 m_a^2$ $\quad V(a_+) = \rho_m \frac{(a_+^2 - a_-)^2}{2\Lambda_a^2}$

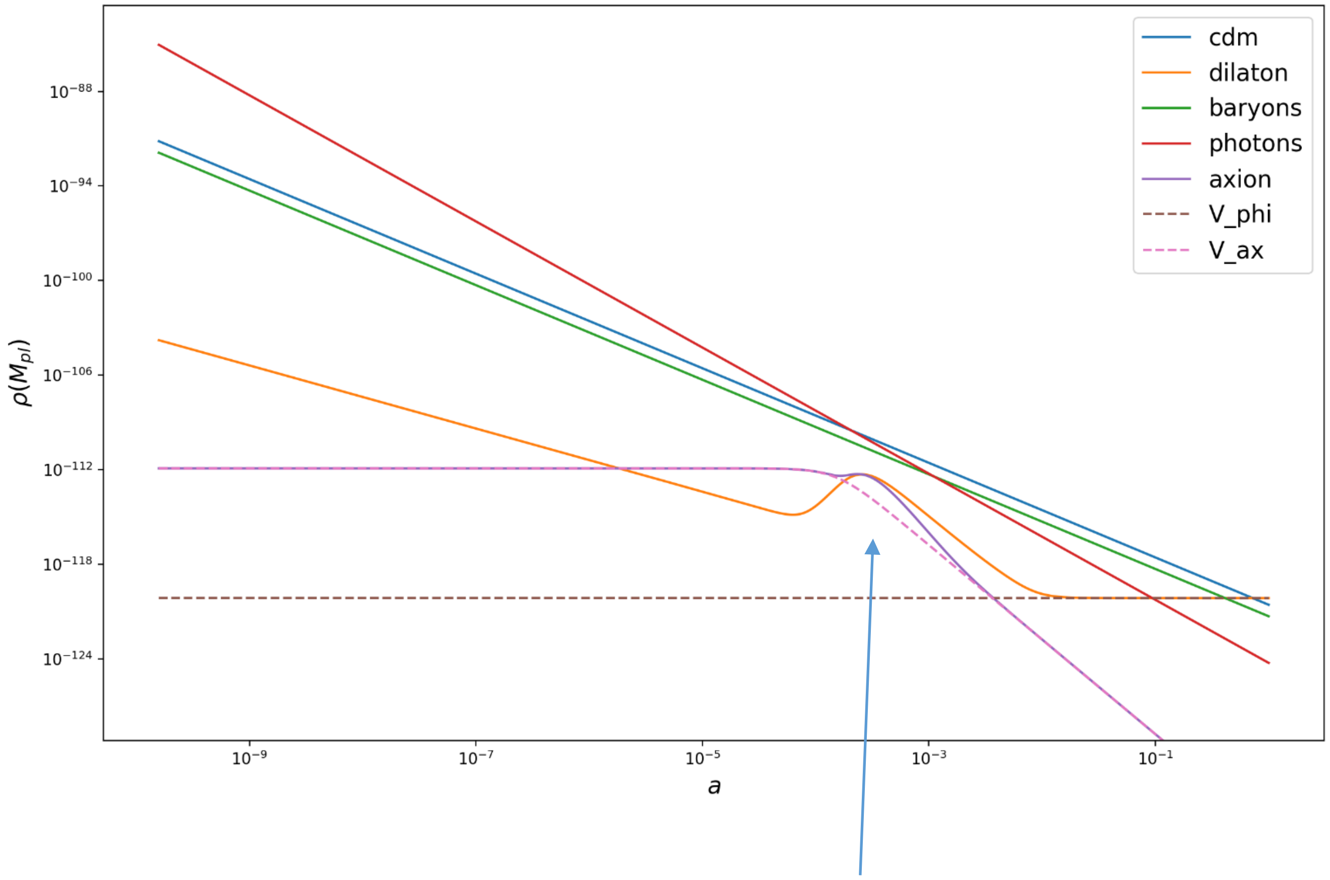
Tachyonic instability:

$$\ddot{\phi} + 3H\dot{\phi} - \dot{a}^2 W W_\phi = -\partial_\phi V_{\text{eff}}(\phi)$$

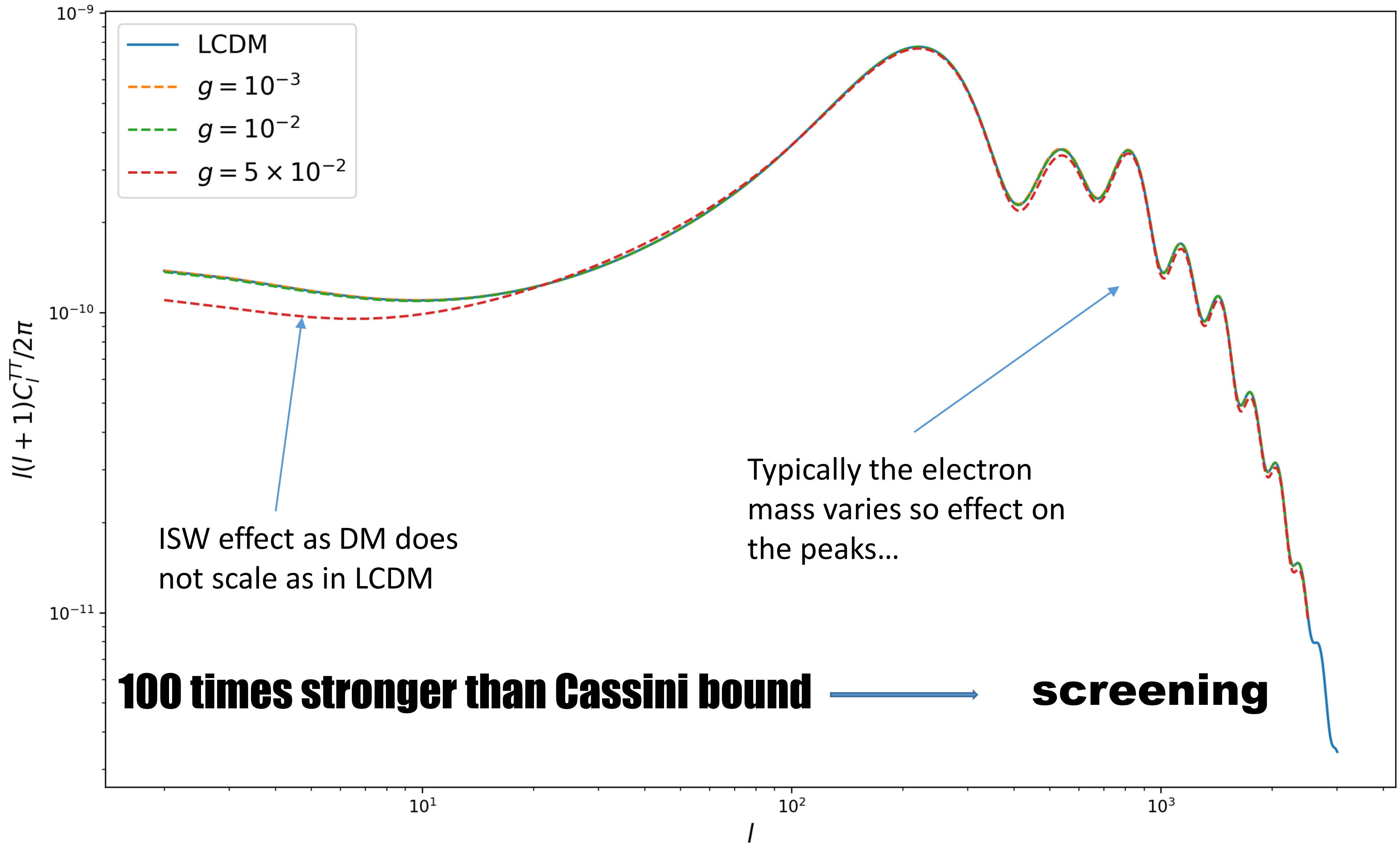
Negative mass term in the quadratic case

No instability for exponential screening.

Exponential potential



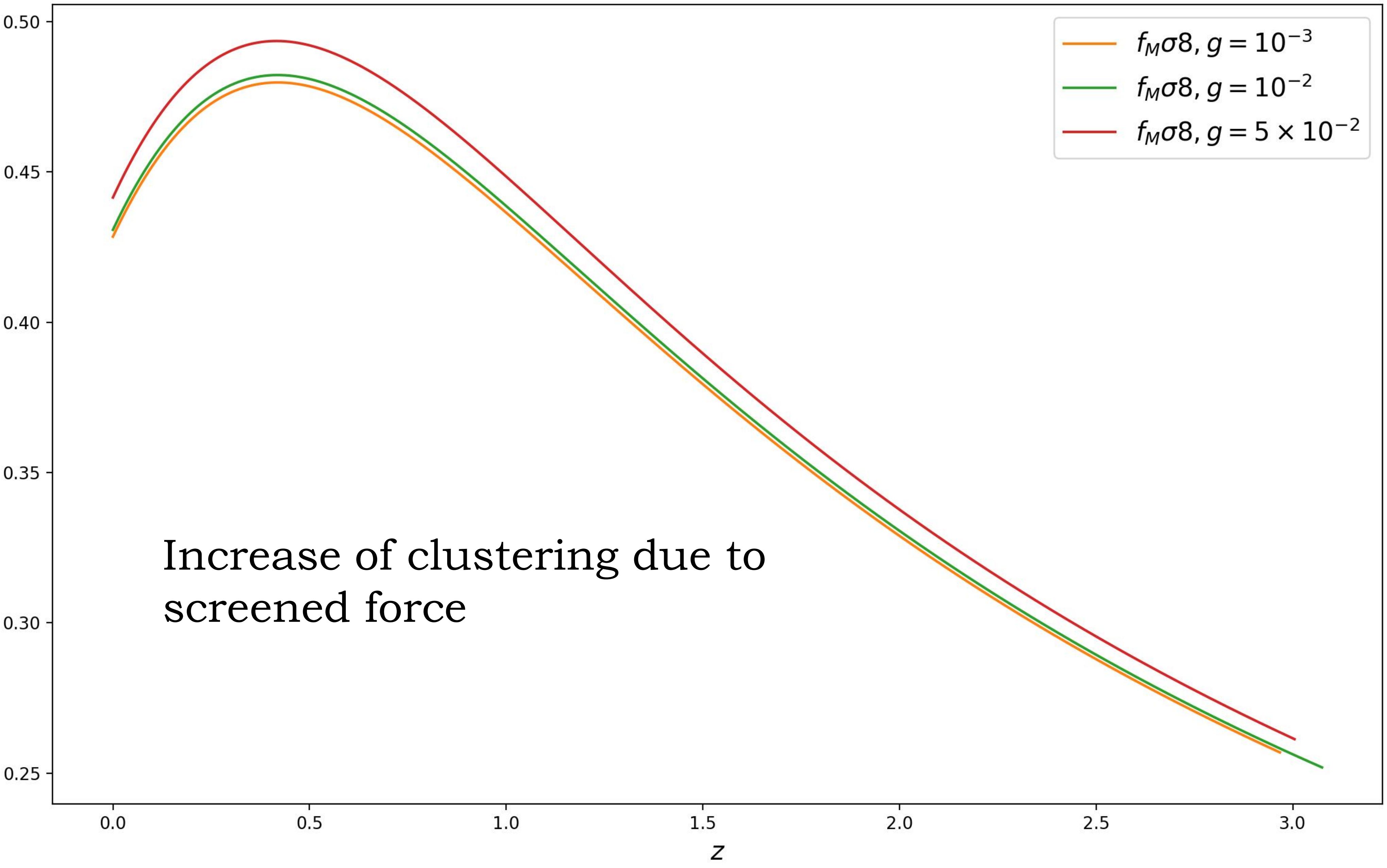
Screening induces an instability limiting the amount of early dark energy



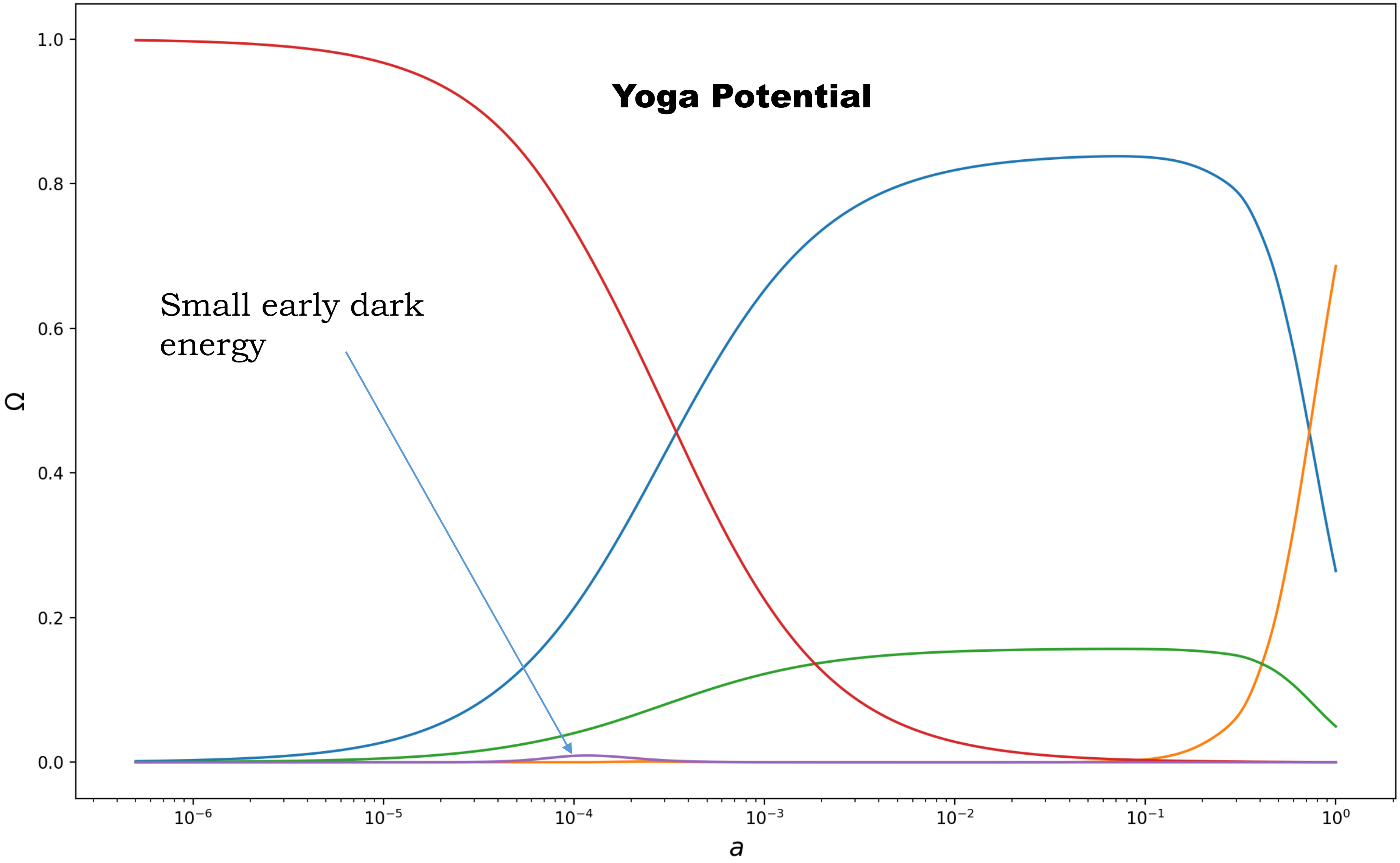
ISW effect as DM does not scale as in LCDM

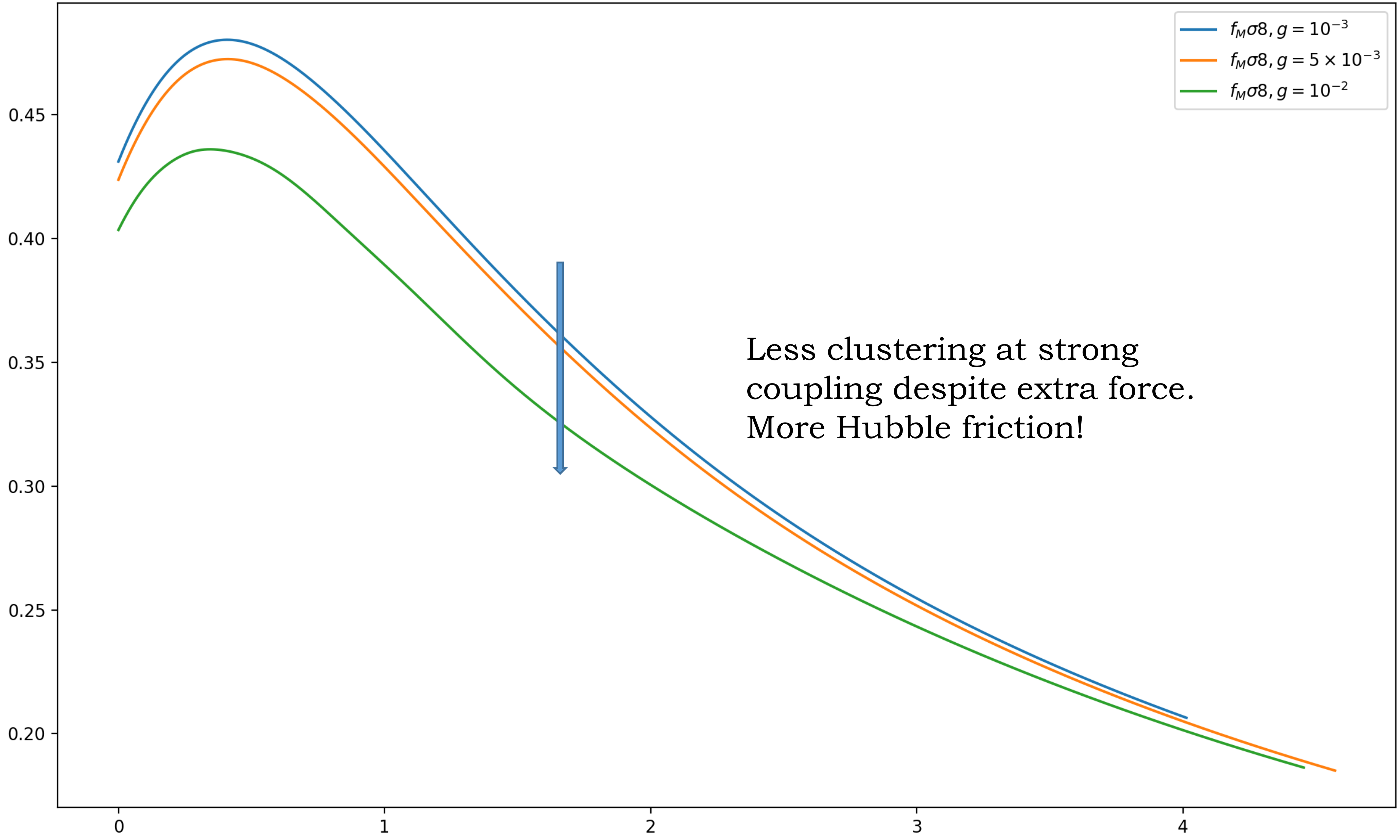
Typically the electron mass varies so effect on the peaks…

100 times stronger than Cassini bound → **screening**



Increase of clustering due to screened force


Yoga Potential
Small early dark energy
Ω
1.0
0.8
0.6
0.4
0.2
0.0
10^{-6}
10^{-5}
10^{-4}
10^{-3}
10^{-2}
10^{-1}
10^{0}
a



Less clustering at strong coupling despite extra force. More Hubble friction!

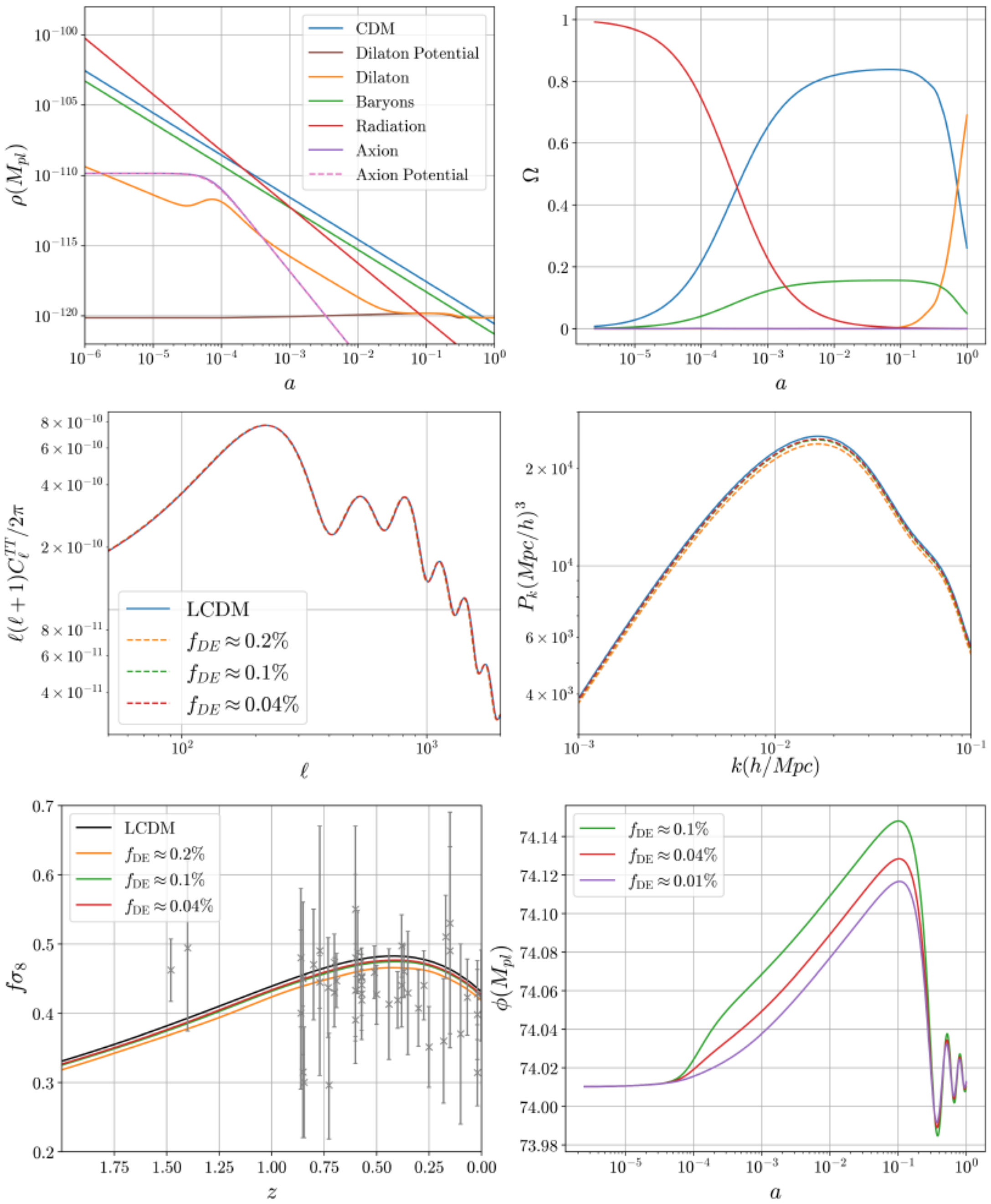


Figure 3: The case of a Yoga potential V for the dilaton and a quadratic W which satisfies screening. Top row shows the background evolution for $m_{\mathfrak{a}} = 5 \times 10^{-14}$eV, $\mathfrak{a}_+ - \mathfrak{a}_- = 2 \times 10^4$GeV, $\Lambda_{\mathfrak{a}} = 10^5$GeV, $\Lambda_\phi = 3 \times 10^3$GeV. Middle row shows the angular and matter power spectra and bottom row shows $f\sigma_8$ and the evolution of the dilaton field for $a_+ - a_- = 2 \times 10^4$GeV, 1.5×10^4GeV, and 1×10^4GeV in orange, green and red respectively. In all plots $\beta = 10^{-2}$.

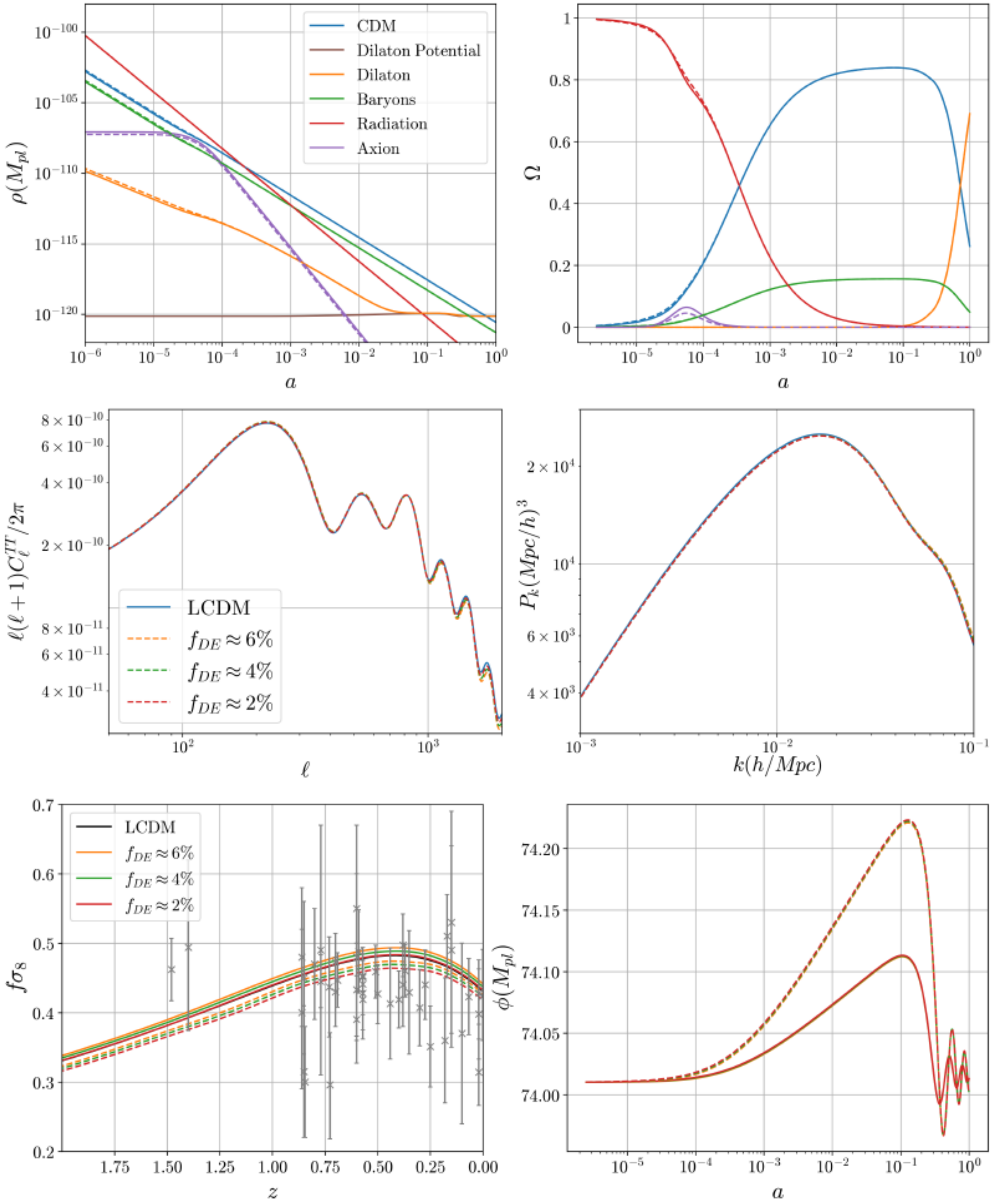


Figure 4: The case of a Yoga potential V and an exponential W which satisfies screening. Top row shows the background evolution for $m_a = 5 \times 10^{-15}$eV, $\Lambda_\mathfrak{a} = 10^6$GeV, $\Lambda_\phi = 2 \times 10^8$GeV, $\mathfrak{a}_+ - \mathfrak{a}_- = 3 \times 10^6$GeV and 4×10^6GeV in solid and dashed respectively. Middle row shows the angular and matter power spectra and bottom row shows $f\sigma_8$ and the evolution of the dilaton field for $\mathfrak{a}_+ - \mathfrak{a}_- = 3 \times 10^6$GeV, 4×10^6GeV and 5×10^6GeV in red, green and orange respectively. The dashed lines in the bottom plots correspond to the same parameters except for $\beta = 2 \times 10^{-2}$. In all other cases, $\beta = 10^{-2}$.

Summary

- Dynamical dark energy would lead to large deviations from General Relativity locally: needs screening.
- Screening can be achieved in the multi-field setting with nothing beyond standard field theory.
- Questions:
  - With one field there are 3 possible screening mechanisms: here with multiple fields??
  - Is there a cosmological signature: ISW? Clustering? Early DE?
  - Are there effects on the equivalence principle: Microscope?
  - Could screening be obtained from the moduli space of a string compactification ?????
- Could we construct a realistic model of dynamical dark energy with screening?????

Contrary to what is usually assumed in cosmology, the scalar models arising from more fundamental (?) theories such as string theory always arise in pairs, and both could play a role:

Supersymmetry → **Superfields=complex fields**

$$\mathcal{L} = K_{a\bar{a}} \partial_\mu z^a \partial^\mu \bar{z}^{\bar{a}}$$

These fields in general have no potential as they arise as moduli in compactifications and one can deform the compactification manifold with no cost (unless there are fluxes…) or non-perturbative phenomena occur (like in QCD).

In general they couple to matter with a conformal coupling

$$g^J_{\mu\nu} = e^{K/3m^2_{\rm Pl}} g_{\mu\nu}$$

Jordan

Einstein

Explicit couplings could also appear for instance the Yukawa couplings could be moduli-dependent

Coupled massless moduli → ☺ **Screening** ☺

??

Burgess and Quevedo concentrated on the axio-dilaton case corresponding to the volume modulus of compactifications from 10d to 4d:

$$K = -3m_{\mathrm{Pl}}^2 \ln(T + \bar{T})$$

The volume modulus can be decomposed in:

$$T = e^{\phi} + ia \quad \Longrightarrow \quad e^{K/6m_{\mathrm{Pl}}^2} = e^{-\phi/2}$$

dilaton

axion

Dilaton coupling to matter of order one!

They also assume that the matter action depends on the axion:

$$\mathcal{A} = -\frac{2}{\sqrt{-g}} \frac{\delta S_m}{\delta a}$$

Energetics

Traditionally for a nearly massless and unscreened field we have two sources of energy:

$$E_{\text{kin}} = 4\pi \int_0^\infty dr r^2 \frac{\phi'^2}{2} = E_0 + \frac{L^2}{4G_N R}$$

Kinetic energy inside independent of the local value of the field

Kinetic energy outside the body also independent of the local value of the field.

$$E_{\text{pot}} = 4\pi \int_0^D dr\ r^2 V(\phi) \simeq \frac{4\pi}{3} D^3 (V_\infty + m_{\text{min}}^2 (\phi_\infty - \phi_{\text{min}})^2) \xrightarrow[\text{Minimised}]{} \phi_\infty = \phi_{\text{min}}$$

IR cutoff

The local field value is the cosmological one locally as long as the local dynamics are subdominant to the cosmological one.

This is only true when

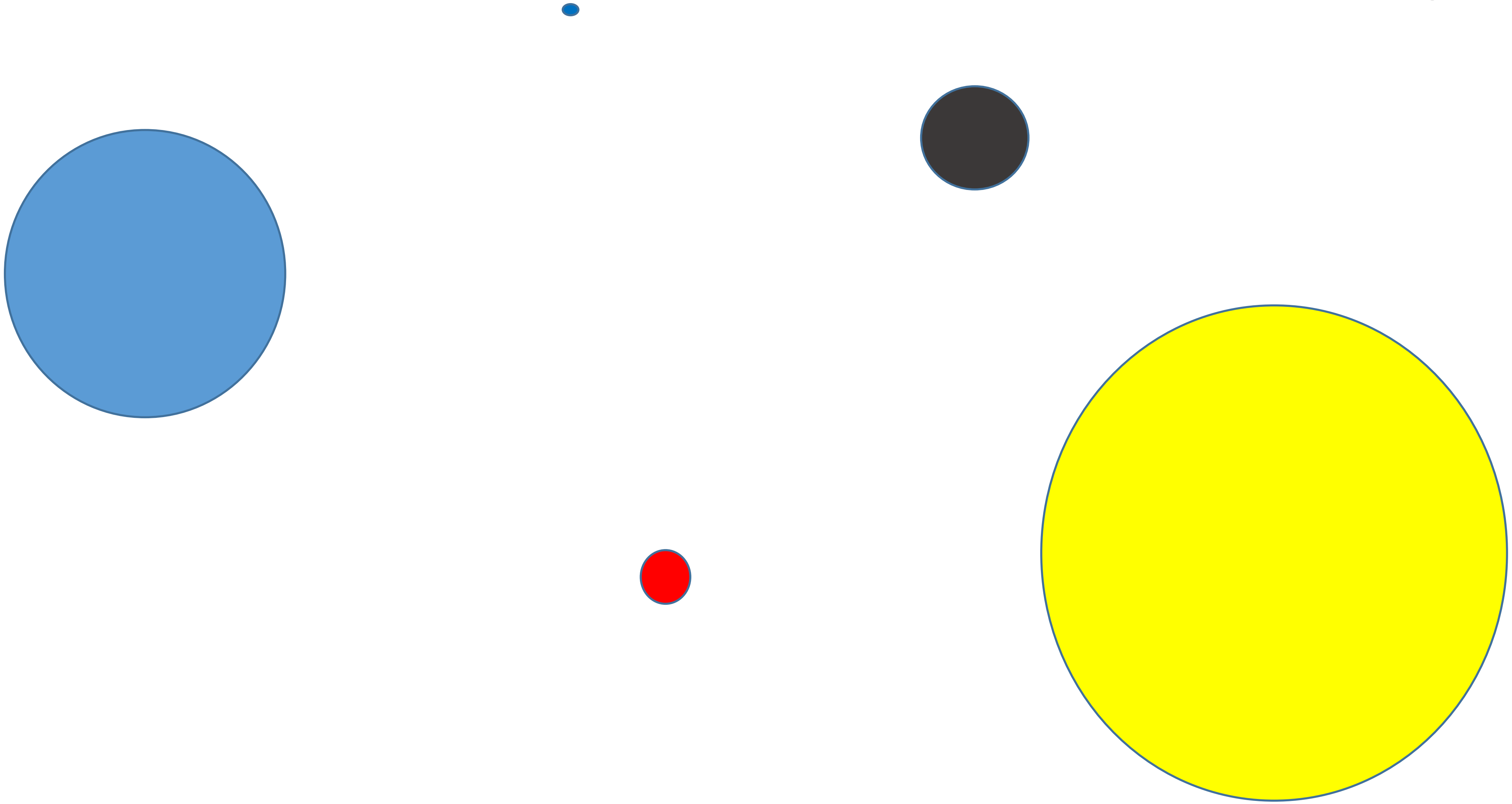
$$(H_0 D)^3 \geq \frac{2\beta^2}{\Omega_\Lambda} \Phi_N^2 (R H_0)$$

For the Sun and the Cassini bound, the cosmological dynamics dominate when D is much larger (3000 a.u) than the solar system.

On small scales, the effects of the local masses dominates and dictates the local value of the field.

Frustration

Who wins?



In the presence of many bodies, close enough to each of them we expect the biggest one to determine the value of the field outside and the screening of all the smaller objects:

$$\frac{\phi_{\rm min} - \phi_\star}{m_{\rm Pl}} \simeq -\frac{8\beta\ell\Lambda_\phi^2}{(a_+ - a_-)^2}\left(\frac{G_N M}{R^2}\right)$$

The object's surface gravity

The object with the **largest surface gravity** wins.

When the *mass of the axion is very large* compared to the inverse size of the object, we can apply the

Narrow Width Approximation

The axion field jumps over a size:

$$\ell \simeq m_a^{-1}$$

$$a = a_+ \Theta(r - R) + a_- \Theta(R - r)$$

$$a' = (a_+ - a_-)\delta(r - R)$$

We will need the expression for the square of a' :

$$(a')^2 = \frac{\delta(r - R)}{2\ell}$$

Same trick as when calculating cross section from the square of the matrix elements .

Minimising the energy with respect to the asymptotic value would determine:

$$W'(r=R)=0$$

If it were not the contribution from the axion surface term in W.

Screening takes place when the scalar is attracted dynamically close to W’=0

$$W^2(\phi)=W_\star^2+\frac{(\phi-\phi_\star)^2}{2\Lambda_\phi^2}+\ldots$$