Unifying inflation with dark energy by healing the cosmological constant problem

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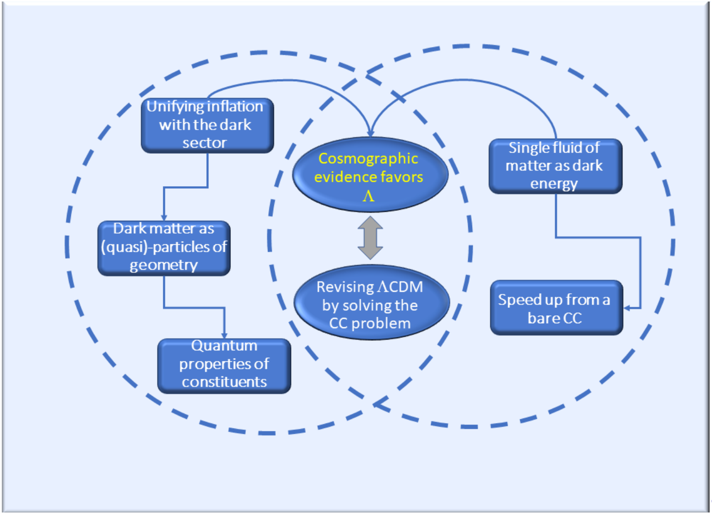
Mind map

No evidence of extending GR from cosmography (and other experimental tests) $\Longrightarrow$ solving the CC problem reconciles inflation with the dark sector $\Longrightarrow$ reformulating the concordance paradigm, namely the ΛCDM model!



Cosmology and Cosmography

The cosmological principle: Homogeneity and isotropy.

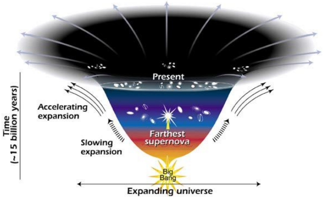
$$ds^2 = dt^2 - a(t)^2 \left[\frac{dr^2}{1-kr^2} + r^2 \left(\sin^2\theta d\phi^2 + d\theta^2 \right) \right]$$

Undisclosed dark energy and dark matter's natures: Understanding their "microphysics" would open new theoretical scenarios.

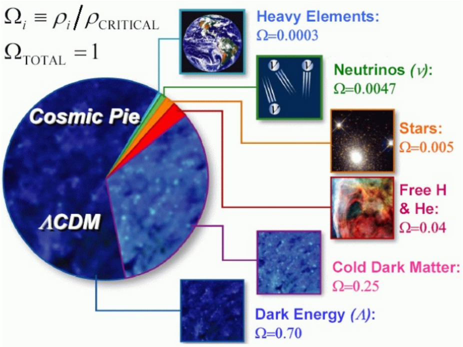
$$H^2 = \frac{1}{3}\rho - \frac{k}{a^2} \qquad \dot{H} + H^2 = \frac{1}{6}(3P + \rho)$$

A model of the Universe should answer these questions:

1. How did the Universe evolve in the past (Big Bang, inflation, etc.)?
2. Is the Universe currently dominated by exotic dark matter and energy?
3. Will the Universe expand forever or will it collapse (what is its destiny)?



The ΛCDM paradigm: The concordance (background) model



...and the concordance model?

The standard ΛCDM paradigm is jeopardized by two main caveats:

- Fine-tuning: $\dfrac{\rho_\Lambda}{\rho_P} \simeq \dfrac{5.96 \times 10^{-27}\ \text{kg/m}^3}{5.16 \times 10^{96}\ \text{kg/m}^3} \sim 10^{-123}$
- Coincidence: $\dfrac{\rho_\Lambda}{\rho_m} \propto a^3$ and yet $\Omega_m \approx 0.3,\ \Omega_\Lambda \approx 0.7$

A model-independent treatment: Cosmography

- Taylor expansion of the scale factor (assuming flat FRW Universe):

$$a(t) = 1 + \sum_{k=1}^{\infty} \frac{1}{k!} \frac{d^k a}{dt^k}\bigg|_{t=t_0} (t - t_0)^k$$

- Cosmographic series:

$$H(t) \equiv \frac{1}{a}\frac{da}{dt}\,, \qquad q(t) \equiv -\frac{1}{aH^2}\frac{d^2a}{dt^2}\,, \qquad j(t) \equiv \frac{1}{aH^3}\frac{d^3a}{dt^3}\,, \qquad s(t) \equiv \frac{1}{aH^4}\frac{d^4a}{dt^4}$$

- Luminosity distance and Hubble expansion rate:

$$d_L(z) = (1+z)\int_0^z \frac{dz'}{H} = \frac{1}{H_0}\left(c_1 z + c_2 z^2 + c_3 z^3 + c_4 z^4\right) + \mathcal{O}(z^5)$$

$$H(z) = \left[\frac{d}{dz}\left(\frac{d_L(z)}{1+z}\right)\right]^{-1} = H_0\left[1 + H^{(1)} z + H^{(2)}\frac{z^2}{2} + H^{(3)}\frac{z^3}{6}\right] + \mathcal{O}(z^4)$$

$$H^{(1)} = 1 + q_0\,,\ H^{(2)} = j_0 - q_0^2\,,\ H^{(3)} = 3q_0^2 + 3q_0^3 - j_0(3 + 4q_0) - s_0$$

[Aviles, Gruber, Luongo, Quevedo, PRD, 86, 123516, (2012)]

- Limits of standard cosmography:
  - the radius of convergence of the Taylor series is restricted to $|z| < 1$;
  - when cosmological data at $z > 1$ are used, the Taylor series do not provide a good approximation of the luminosity distance due to its divergent behavior;
  - finite truncations propagate errors that may result in possible misleading outcomes.

- Introducing a "new" cosmography: The use of rational polynomials
  - they extend the radius of convergence of Taylor series;
  - they can better approximate situations at high-redshift domains;
  - the series can be modelled by choosing appropriate orders depending on each case of interest.

$\Longrightarrow$ Proposing the first attempts toward high redshift cosmography:

1. Padé rational series.
2. Chebyshev rational polynomials.

[Dunsby, Luongo, IJGMMP, 13, 1630002 (2016)]

Cosmography with Padé and Chebyshev polynomials

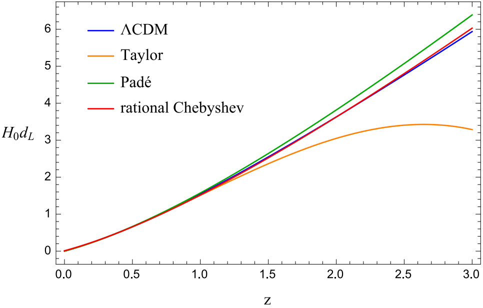
- **Padé polynomials:** Given $f(z)=\sum_{i=0}^{\infty} c_i z^i$, $c_i = f^{(i)}(0)/i!$, the (N,M) Padé series reads

$$P_{N,M}(z)=\frac{\sum_{n=0}^{N} a_n z^n}{1+\sum_{m=1}^{M} b_m z^m}\,, \qquad \begin{cases} P_{N,M}(0)=f(0) \\ P'_{N,M}(0)=f'(0) \\ \vdots \\ P_{N,M}^{(N+M)}(0)=f^{(N+M)}(0) \end{cases}$$

- **Chebyshev polynomials:** Given the first kind Chebyshev polynomials $T_n(z)=\cos(n\theta)$, $n\in\mathbb{N}_0$, $\theta=\arccos(z)$, with $f(z)=\sum_{k=0}^{\infty}{}' c_k T_k(z)$, a (n,m) rational Chebyshev approximation of $f(z)$ reads:

$$R_{n,m}(z)=\frac{\sum_{i=0}^{n}{}'\ a_i T_i(z)}{\sum_{j=0}^{m}{}'\ b_j T_j(z)}$$

[Gruber, Luongo, PRD, 89, 103506 (2014)]



Comparison among different cosmographic techniques

Figure: (2,1) rational Chebyshev approximation of the luminosity distance for the ΛCDM model with the correspondent Padé and Taylor approximations.

Parameter	Taylor			Padé			Rational Chebyshev		
	Mean	1σ	R.E.	Mean	1σ	R.E.	Mean	1σ	R.E.
H_0	65.80	$^{+2.09}_{-2.11}$	3.19%	64.94	$^{+2.11}_{-2.02}$	3.17%	64.95	$^{+1.89}_{-1.94}$	2.95%
q_0	-0.276	$^{+0.043}_{-0.049}$	16.8%	-0.285	$^{+0.040}_{-0.046}$	15.1%	-0.278	$^{+0.021}_{-0.021}$	7.66%
j_0	-0.023	$^{+0.317}_{-0.397}$	1534%	0.545	$^{+0.463}_{-0.652}$	102%	1.585	$^{+0.497}_{-0.914}$	44.5%

Table: 68% confidence level constraints and relative errors from the MCMC analysis of SN+OHD+BAO data for the fourth-order Taylor, (2,2) Padé and (2,1) rational Chebyshev polynomial approximations of the luminosity distance.

[Capozziello, D'Agostino, Luongo, MNRAS, 476, 3924 (2018)]

Modified gravity: $f(T)$ and Extended gravity: $f(R)$

- $f(T)$ **action:**

$$S=\int d^4x\; e\left[\frac{f(T)}{2}+\mathcal{L}_m\right],\qquad e=\sqrt{-g}=\det(e^A_\mu)$$

- Modified Friedmann equations:

$$H^2=\frac{1}{3}(\rho_m+\rho_T),\quad 2\dot{H}+3H^2=-\frac{1}{3}(p_m+p_T)$$

with $\rho_T=Tf'(T)-\dfrac{f(T)}{2}-\dfrac{T}{2}\,,\qquad p_T=\dfrac{f-Tf'(T)+2T^2f''(T)}{2[f'(T)+2Tf''(T)]}.$

- $f(R)$ **action:**

$$S=\int d^4x\sqrt{-g}\left[\frac{f(R)}{2}+\mathcal{L}_m\right]$$

- Extended Friedmann equations:

$$H^2=\frac{1}{3}\left[\frac{1}{f'}\rho_m+\rho_{curv}\right]\quad 2\dot{H}+3H^2=-\frac{p_m}{f'}-p_{curv}$$

with $\rho_{curv}=\dfrac{1}{2f'}(f-Rf')-3H\dot{R}\frac{f''}{f'},\ p_{curv}=2H\dot{R}\frac{f''}{f'}+\ddot{R}\frac{f''}{f'}+\dot{R}^2\frac{f'''}{f'}-\frac{1}{2f'}(f-Rf')$

[D'Agostino, Luongo, PRD, 98, 124013 (2018)]

[Abedi, Capozziello, D'Agostino, Luongo, PRD, 97, 084008 (2018)]

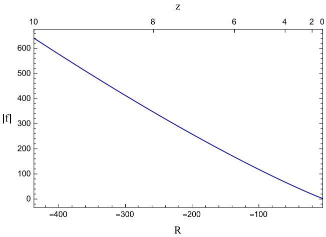
Reconstruction of the $f(T)$ and $f(R)$ actions

Common strategy: Combining the Friedmann equations with initial conditions:

1. $G_{\rm eff} \equiv G_N/f'(T) \Longrightarrow \left.\frac{df}{dz}\right|_{z=0} = 1,\ f(T(z=0)) = f_0 = 6{H_0}^2(\Omega_{m0} - 2)$
2. $G_{\rm eff} = G_N/f'(R) \Longrightarrow f'(R_0) = 1$ with $f_0 = R_0 + 6H_0^2(\Omega_{m0} - 1)$

Moreover, considering the cosmographic $H_{2,1}(z)$ then:

Use $T = -6H^2$ to find $z(T)$ and plug $z(T)$ into $f(z)$ to find $f(T)$.
Use $R = -6(\dot{H} + 2H^2)$ to get $z(R)$ and plug into $f(z)$ to obtain $f(R)$.



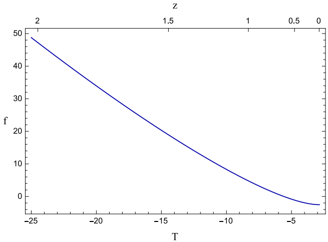


Figure: Numerical reconstruction of $f(T)$ and $f(R)$ through (2,1) Padé approximation of $H(z)$.
[Aviles, Bravetti, Capozziello, Luongo, PRD, 87, 064025, (2013) & PRD, 90, 044016, (2014).]
[Capozziello, D'Agostino, Luongo, GeRG, 49, 141, (2017) & JCAP, 1805, 008, (2018).]

A single matter-like fluid with pressure?

The effective representation of *dust with pressure* in a curved space-time is given by:

$$\mathcal{L}_1 = K(X,\varphi) + \lambda Y[X,\nu(\varphi)]\ , \qquad (1)$$

$$\mathcal{L}_2 = -V^{\text{eff}}(X,\varphi)\ . \qquad (2)$$

We require:

- There exists only *one fluid*, composed of BM and DM and *vacuum energy*.
- Matter is coupled to Λ through $V^{\text{eff}}(X,\varphi)$.
- The coupling should cancel vacuum energy through a *first order phase transition*.
- Thermodynamics *naturally* leads to a negative pressure.
- The phase *during* the phase transition produces (quasi)-particles of dark matter.

[Luongo, Quevedo, AstSpS, 338, 2, 345-349, (2012)]
[Luongo, Quevedo, GeRG, 46, 1649, (2014)]
[Dunsby, Luongo, Reverberi, PRD, 94, 083525, (2016)]
[Luongo, Muccino, PRD, 98, 103520, (2018)]
[Belfiglio, Luongo, Mancini, PRD, 105, 123523, (2022)]
[Belfiglio, Giambò, Luongo, ArXiv:2206.14158, Submitted to PRD, (2022)]

A single matter-like fluid with pressure?

Our fluid consists of BM and DM, so:

$$\mathcal{L}_1 = K_{\rm BM} + K_{\rm DM} + \lambda\left(Y_{\rm BM} + Y_{\rm DM}\right)\,, \tag{3}$$

The Lagrangian $\mathcal{L}_2$ models the coupling with quantum vacuum energy.

In a thermal bath $\rightarrow$ **spontaneous symmetry breaking!**

$$V^{\rm eff}(X,\varphi) = V_0 + \frac{\chi}{4}\left(\varphi^2 - \varphi_0^2\right)^2 + \frac{\chi}{2}\varphi_0^2\varphi^2\,\frac{T^2(X)}{T_{\rm c}^2}\,, \tag{4}$$

where $T_{\rm c} = \varphi_0\sqrt{\chi/\tilde{g}}$ is the critical temperature.

BT, $T > T_{\rm c}$: the minimum is at $\varphi = 0$ and its value is $V_0 + \chi\varphi_0^4/4$.
AT, $T < T_{\rm c}$: the minimum is at $\varphi = \varphi_0$ and its value is V_0.

As $T_{\alpha\beta} = 2X\mathcal{L}_X v_\alpha v_\beta - \left(K - V^{\rm eff}\right) g_{\alpha\beta} \implies$ density and pressure are:
$\rho\left(\lambda, X, \varphi\right) = 2X\mathcal{L}_X - \left(K - V^{\rm eff}\right)$ *and* $P\left(X,\varphi\right) = K - V^{\rm eff}$.

$Y = 0$ varying w.r.t. λ $\qquad$ $\mathcal{L}_\varphi - \nabla_\alpha\left(\mathcal{L}_X \nabla^\alpha \varphi\right) = 0$ varying w.r.t. φ

Thermodynamics

Non-dissipative fluids are described by virtue of the *pullback* formalism through Carter's covariant formulation. Fluids are framed with four scalar fields, namely ϕ^a, $a = 1, 2, 3$. Fluids evolve as comoving coordinates: ϕ^0 is an internal time coordinate.

These scalars can be viewed as Stückelberg fields.

Consequently the shift symmetry is valid

$$\boxed{\varphi \to \varphi + c^0}$$

First principle, Gibbs-Duhem relation and Helmotz free-energy density:
$d\rho = Tds + \mu dn, \quad dP = sdT + nd\mu, \quad df = \mu dn - sdT$

Combining all together:

$$f = -\mathcal{L}, \qquad s = \sqrt{2X}\mathcal{L}_X, \qquad T = \sqrt{2X}, \qquad \mu = 0$$

fulfilling the conditions:

$$\left.\frac{\partial\,(fV)}{\partial T}\right|_V = -\sqrt{2X}\mathcal{L}_X V = -sV\,, \left.\frac{\partial\,(fV)}{\partial V}\right|_T = -\mathcal{L} = -\left(K - V^{\rm eff}\right) = -P\,.$$

The above relations must be consistent with standard thermodynamics.

This naturally provides the sign of P at all times: $f > 0 \to P < 0$

Noether current and entropy

Dust-like matter having pressure **naturally fixes the sign of P to be negative**. This ensures no need of putting by hand the sign of P inside Einstein's equations.

Noether's theorem

The global shift symmetry changes the matter Lagrangian density $\mathcal{L}_1$ mostly by a total divergence.

We explicitly get:

$$\begin{aligned}
&\mathcal{L}_1\left(X', \varphi'\right) = \mathcal{L}_1\left(\frac{1}{2}\nabla_\alpha\varphi\nabla^\alpha\varphi, \varphi + c^0\right) = \\
&\mathcal{L}_1\left(X, \varphi\right) + c^0\left[\frac{\partial\mathcal{L}_1}{\partial\varphi} - \nabla_\alpha\frac{\partial\mathcal{L}_1}{\partial\left(\nabla_\alpha\varphi\right)}\right] + c^0\nabla_\alpha\frac{\partial\mathcal{L}_1}{\partial\left(\nabla_\alpha\varphi\right)} = \\
&\mathcal{L}_1\left(X, \varphi\right) + c^0\nabla_\alpha\left(\mathcal{L}_{1,X}\nabla_\alpha\varphi\right)\,,
\end{aligned} \tag{5}$$

where, in the second line of Eq. (5), the quantity in the brackets identically vanishes in view of the Euler–Lagrange equations.

The conserved current $\mathcal{J}_1^\alpha$ corresponds to the total divergence of Eq. (5). So $\mathcal{J}_1^\alpha = \sqrt{2X}\left(K_X + \lambda Y_X\right) v^\alpha$

Noether current and entropy

How do we deal with $V^{\rm eff}$?

before the transition (**BT**, with $V^{\rm eff} = V_0 + \chi\varphi_0^4/4$ at $\varphi = 0$)
after the transition (**AT**, with $V^{\rm eff} = V_0$ at $\varphi = \varphi_0$)

During the two phases, the Noether's theorem implies that:

$$\mathcal{L}_2\left(X'\right) = -V^{\rm eff}(X) - c^0\nabla_\alpha\left(V_X^{\rm eff}\nabla_\alpha\varphi\right), \tag{6}$$

where another conserved current from the total divergence of Eq. (6): $\mathcal{J}_2^\alpha = -\sqrt{2X}V_X^{\rm eff}v^\alpha$, so the total conserved current $\mathcal{J}^\alpha$ is the entropy density current $s_\alpha = sv_\alpha$, namely $\mathcal{J}^\alpha \equiv \mathcal{J}_1^\alpha + \mathcal{J}_2^\alpha = \sqrt{2X}\mathcal{L}_X v^\alpha = s^\alpha$. Conserved currents imply

$\mathcal{L}_\varphi = 0$: *The Lagrangian does not depend upon* φ.
$\mathcal{L}(\lambda, X, \nu) = K(X) - V^{\rm eff}(X) + \lambda Y(X, \nu)$

We simply recover the standard Euler relation $P + \rho = Ts + \mu n$ and recast the energy-momentum tensor as: $T_{\alpha\beta} = (Ts_\alpha + \mu n_\alpha)v_\beta + Pg_{\alpha\beta}$, where $n_\alpha = nv_\alpha$ is the particle number density current.

The projection of the energy-momentum tensor conservation along v^α, i.e., $v^\alpha\nabla^\beta T_{\alpha\beta} = 0$, leads to $T\nabla^\alpha s_\alpha + \mu\nabla^\alpha n_\alpha = 0$. So by virtue of the existence of $\mathcal{J}^\alpha$, it reduces to $\mu\nabla^\alpha n_\alpha = 0$, but also $\mu = 0$! $\Longrightarrow$ No particles produced, **BT** and **AT**!

Properties of the fluid: Small perturbations

Main consequences

The conservation of the energy-momentum tensor can be recast by Carter-Lichnerowicz equations

$$n\mathcal{W}_{\alpha\nu} v^\nu = nT\nabla_\alpha\sigma - \varsigma_\alpha \nabla^\nu n_\nu$$

where $\mathcal{W}_{\alpha\nu} = \nabla_\nu\varsigma_\alpha - \nabla_\alpha\varsigma_\nu$ is the vorticity tensor, $\varsigma^\alpha = h/nv^\alpha$ the current of the enthalpy per particle, and $\sigma = s/n$ the entropy per particle. We immediately infer:

$$\mathcal{W}_{\alpha\nu} = 0 \quad \Rightarrow \quad \text{the fluid is irrotational}$$
$$\nabla_\alpha\sigma = 0 \quad \Rightarrow \quad \text{the fluid is isentropic}$$

Small perturbations

Conformal Newtonian gauge: $ds^2 = a(\tau)^2\left[(1+2\Phi)\,d\tau^2 - (1-2\Phi)\,dx^2\right]$, with Φ is the Newtonian potential, $\tau = a(t)t$ the conformal time. The entropy perturbation shift, Δ, reads:

$$\Delta = \left(\frac{\delta P}{\delta\rho} - c_{\rm s}^2\right)\frac{\delta\rho}{P} = -\frac{D(X)\delta\nu + E(X)\delta\lambda}{P}$$

Requiring $Y_X \neq 0$, we find:

1. $c_{\rm s}^2 \equiv 0$ and $P = \text{const}$
2. Minimum of Gibbs energy $\Longrightarrow$ fluid at equilibrium.
3. $P = \text{const} \Longrightarrow T = \text{const}$ *in the proximity of each minima of* $V^{\rm eff}$.

What about quantum vacuum energy?

The cosmological constant problem $\Lambda = \rho_{\rm vac} + \Lambda_B$

Cancellation mechanism to remove $\rho_{\rm vac} \implies \Lambda_B = \rho_\Lambda^{\rm obs} = \Lambda$ today!

To solve this problem: What's the role of $V^{\rm eff}$? **We thus explore two possibilities:**

1) $V_0 = -\chi\varphi_0^4/4$, so BT we have $V^{\rm eff} = 0$ and hence

$$P_1 = \begin{cases} K & \rm (BT) \\ K + \chi\varphi_0^4/4 & \rm (AT) \end{cases}, \tag{7}$$

$$\rho_1 = \begin{cases} 2X\lambda Y_X - K & \rm (BT) \\ 2X\lambda Y_X - K - \chi\varphi_0^4/4 & \rm (AT) \end{cases}, \tag{8}$$

and since $f > 0$, then $K < -\chi\varphi_0^4/4$.

2) $V_0 = 0$, so AT we have $V^{\rm eff} = 0$ and hence

$$P_2 = \begin{cases} K - \chi\varphi_0^4/4 & \rm (BT) \\ K & \rm (AT) \end{cases}, \tag{9}$$

$$\rho_2 = \begin{cases} 2X\lambda Y_X - K + \chi\varphi_0^4/4 & \rm (BT) \\ 2X\lambda Y_X - K & \rm (AT) \end{cases}, \tag{10}$$

and since $f > 0$, then $K < 0$.

In both cases $K < 0$, but with different magnitudes!

Focusing on the two cases

Case $V_0 = 0$.

$$P_2 \approx \begin{cases} -\chi\varphi_0^4/4 & \text{(BT)} \\ K & \text{(AT)} \end{cases}, \tag{11}$$

$$\rho_2 = \begin{cases} (\rho_{\rm DM} + \rho_{\rm BM})(1+z)^3 + \chi\varphi_0^4/4 & \text{(BT)} \\ (\rho_{\rm DM} + \rho_{\rm BM})(1+z)^3 - K & \text{(AT)} \end{cases}, \tag{12}$$

where the BM can be considered even pressureless.

Vacuum energy density cancels without dark matter $\Longrightarrow$ discontinuity of $V^{\rm eff}$ $\Longrightarrow$ **This case still suffers from coincidence as the ΛCDM model!**

Case $V_0 = -\chi\varphi_0^4/4$

$$P_1 \approx \begin{cases} K_{\rm DM} & \text{(BT)} \\ K_{\rm BM} & \text{(AT)} \end{cases}, \tag{13}$$

$$\rho_1 \approx \begin{cases} (\rho_{\rm DM} + \rho_{\rm BM})(1+z)^3 - K_{\rm DM} & \text{(BT)} \\ (\rho_{\rm DM} + \rho_{\rm BM})(1+z)^3 - K_{\rm BM} & \text{(AT)} \end{cases}, \tag{14}$$

where $\rho_{\rm BM} = 2X\lambda_0 Y_{{\rm BM},X}$ and $\rho_{\rm DM} = 2X\lambda_0 Y_{{\rm DM},X}$ are constants. **A few consequences:**

Since $K_{\rm DM} \approx -\chi\varphi_0^4/4$ and $K_{\rm BM} \ll K_{\rm DM}$, vacuum energy is elided by dark matter.
As $\chi > 0$, the sign of $K_{\rm DM}$ is opposite to the vacuum energy term.
AT the Universe accelerates *because of a $\Lambda_B \approx -K_{\rm BM}$ contribution*.
Coincidence and fine-tuning: Solved!

But... what about Λ_B and cancelled density?

During the transition: Scalar field (inflaton) Lagrangian

$$\mathcal{L} = \frac{1}{2}\left[g^{\mu\nu}\phi_{,\mu}\phi_{,\nu} - \xi R\phi^2 - 2V(\phi)\right]$$

$\Longrightarrow$ Vacuum energy drives inflation, $\Lambda^4 \sim \rho_{vac} \sim 10^{64}$ GeV

During inflation: $a(\tau) = 1/(1 - H_I\tau)$ and conformal coupling, $\xi = 1/6$. Thus,

- $g_{\mu\nu} = a^2(\tau)\left(\eta_{\mu\nu} + h_{\mu\nu}(x)\right)$ with $h_{\mu\nu}(x)$ *inflaton perturbations*.
- Quadratic hilltop potential, $V(\phi) = \Lambda^4(1 - \phi^2/\mu_2^2)$.

At first order in perturbations:

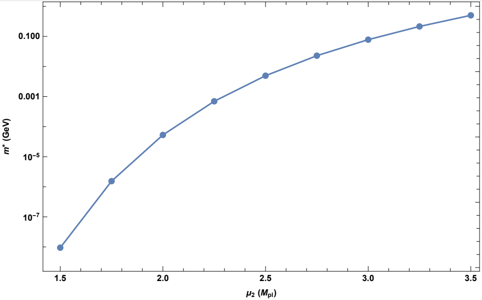
- Interacting Lagrangian: $\mathcal{L}_I = -\frac{1}{2}H^{\mu\nu}T^{(0)}_{\mu\nu}$, with $H_{\mu\nu} = a^2(\tau)h_{\mu\nu}$ and $T^{(0)}_{\mu\nu}$ zero-order energy-momentum tensor.
- The scattering matrix, $\hat{S} \simeq 1 + i\hat{T}\int d^4x\,\sqrt{g_{(0)}}\mathcal{L}_I$.

$\Longrightarrow$ Interaction between field and geometry creates (quasi)-particles of geometry!

$$N^{(2)}(\tau^*) = \frac{a^{-3}(\tau^*)}{(2\pi)^3}\int d^3k\;d^3p\;|\langle 0|\hat{S}|k,p\rangle|^2$$

Interpreting (quasi)-particles as geometric dark matter

No couplings of the inflaton field with standard model fields are involved.
$\Longrightarrow$ We could interpret geometric particles in terms of dark matter candidates, arising from field-curvature coupling.



Predictions of dark matter candidate!

Figure: Mass m^* of the dark matter candidate as function of the hilltop parameter μ_2.
Setting $\Lambda^4 = 10^{64}$, the predicted mass range is $10^{-7} \leq m^\star \leq 10^{-1} GeV$

Vacuum energy cancellation: At the transition ($\tau = 0$) between inflation and radiation/matter epoch there is a pressure shift: $\Delta P \propto -H_I^2 \Longrightarrow$
This mimics a dark fluid with constant ΔP dominating at our times and:

1. If $H_I^2 \equiv \Lambda_B$ the cosmological constant problem is solved!
2. The Israel-Darmois junction conditions hold at $\tau = 0$.

Geometric cosmological entanglement

Perturbations also *create entanglement* in the final state of the system of particles

$$|\Psi\rangle_{\rm in} = |0_k;0_p\rangle_{\rm in} \longrightarrow |\Psi\rangle_{\rm out} = \mathcal{N}\left(|0_k;0_p\rangle_{\rm in} + \frac{1}{2}\hat{S}^{(1)}_{kp}|1_k;1_p\rangle_{\rm in}\right)$$

Entanglement can be quantified using *von Neumann entropy* $\mathcal{S}$ of the reduced density operator

$$\rho_k = \mathrm{Tr}_p\left(|\Psi\rangle_{\rm out}\langle\Psi|\right) \implies \mathcal{S}(\rho_k) = -\mathrm{Tr}_k\left(\rho_k \log_2 \rho_k\right) \neq 0$$

Geometric (quasi)-particles of dark matter do not interact with other particles. So the generated entanglement may be preserved to our time!

This may lead to:

- Entanglement extraction.
- Deduction of cosmological parameters.
- Characterize dark matter nature.
- Unify the dark sector.

Summary and perspectives

...Talk mainly based on:

Cosmography at low and high redshifts: No strong evidence for extending the ΛCDM model!

One dark fluid: Matter-like with pressure to resolve the cosmological constant problem.

Unifying the inflationary scenario: Predicting (quasi)-particles of geometry as dark matter.

...Perspectives and future developments mainly based on:

Testing new model-independent techniques: Cosmography at high redshifts with GRBs, σ_8, GW.

Accretion disks in galaxies with (quasi)-particles of dark matter. Consequences in spiral galaxies.

Resolving the cosmological tensions adopting the here-proposed model.

Quantum information in geometric dark matter framework.