The Weak Scale as a Trigger

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UV landscape \[ \Lambda_{UV} = \frac{M^4}{N_{UV}} \sim v^4 \]

IR landscape \[ \Lambda_{IR} = \frac{\Lambda_{UV}}{N_{IR}} \sim \frac{v^8}{M^4} \]
\[ N_{IR} = 2^{n_{\phi}} \]
\[ m_{\phi} \lesssim \frac{v^2}{M_*} \]

Type 0 2HDM

Light scalar dark matter from EWPT
Question #1
What varies as we change the Higgs mass parameter in the SM?

Answer to Q#1
All the spectrum of the Standard Model including W and Z bosons, quarks and leptons and the Higgs boson itself.

Question #2
Is there any gauge invariant local operator which has a value sensitive to the Higgs mass parameter?
Answer to Q#2

\[ \mathcal{O}_h = h^\dagger h \]

However, \( \mathcal{O}_h \) is not calculable in the SM

Probing by \( \xi \phi h^\dagger h \)

1 loop tadpole is generated

\[ \frac{1}{16\pi^2} \xi \phi \Lambda_H^2 \]

\( \Lambda_H \): cutoff of the Higgs loop

\( \langle h^\dagger h \rangle \) is independent of \( m_h^2 \)
depends on \( \Lambda_H^2 \)
Hierarchy Problem

making $m_h^2$ to be calculable

Closely related question

Is $\langle h^\dagger h \rangle$ calculable?

Two calculable examples

Supersymmetry

$\langle h^\dagger h \rangle \sim m_{\text{SUSY}}^2$

Composite Higgs

$\langle h^\dagger h \rangle \sim f_\pi^2$
$O_G = \text{tr} G \tilde{G}$

SM
Possible operators?

\[ \mathcal{O}_q = qhu^c \]

\[ y_q \mathcal{O}_q + \xi \phi \mathcal{O}_q + \text{h.c.} \]

\[ \frac{\xi y_q^*}{(16\pi^2)^2} \phi \Lambda^4 \]

\[ \phi : \text{dimensionless} \]

\[ \frac{\phi}{M_*} \]

Massless up quark provides the operator \( \mathcal{O}_u = qhu^c \)

one of the solutions to the strong CP problem but is not viable any longer
\[ \mathcal{O}_G = \text{tr} G \tilde{G} \]
\[ \mathcal{O}_G = \partial_\mu K^\mu \]

\[ \langle G \tilde{G} \rangle \sim \theta (m_u + m_d) \Lambda_{QCD}^3 \]

depends on the weak scale and is insensitive to UV

size is too small \[ \Lambda_* \sim (100 \text{ keV})^4 \] strong CP
$O_H = H_1 H_2$

2HDM
Weak scale as a trigger

Type 0 2HDM \( \rightarrow B\mu \) is forbidden

We need a symmetry under which the operator is charged. Otherwise, the operator is UV sensitive (e.g., Yukawa term)

\[
\phi \rightarrow -\phi \\
H_1 H_2 \rightarrow -H_1 H_2
\]

\( Z_4 \) symmetry

<table>
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<tr>
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<th>PQ</th>
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<th>( Z_2 )</th>
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<tr>
<td>( H_1 H_2 )</td>
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<td>( B\mu )</td>
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<td>( \lambda_6 )</td>
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<td>( \lambda_7 )</td>
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<td>( \lambda_5 )</td>
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\( H_1 \rightarrow +iH_1, H_2 \rightarrow +iH_2, \)

\( (H_1 H_2) \rightarrow -(H_1 H_2), \)

\( (q u^c) \rightarrow +i(q u^c), \)

\( (q d^c) \rightarrow -i(q d^c), \)

\( (l e^c) \rightarrow -i(l e^c). \)
\(-\mathcal{L} \supset m_{H_2}^2 |H_2|^2 + m_{H_1}^2 |H_1|^2 + \frac{\lambda_1}{2} |H_1|^4 + \frac{\lambda_2}{2} |H_2|^4 + \lambda_3 |H_2|^2 |H_1|^2 + \lambda_4 |H_1 H_2|^2 + \left( \frac{\lambda_5}{2} (H_1 H_2)^2 + \text{h.c.} \right) \)

Peccei-Quinn symmetry is explicitly broken

\( B\mu \) generates tadpole and should be forbidden
Weak scale as a trigger

\[ \mu^2 \equiv \langle \mathcal{O}_H \rangle \equiv \langle H_1 H_2 \rangle \]

as a function of \( m_h^2 \)

\( \Lambda \) can be small only if \( \mu_S^2 \leq \mu^2 \leq \mu_B^2 \)
Values of $\mu^2$ in the landscape (classical)

(quartic couplings are taken to be order one)

\[ \langle H_1 H_2 \rangle = 0 \]

\[ \langle H_1 H_2 \rangle = \sqrt{|m_{H_1}^2||m_{H_2}^2|} \]

\[ \langle H_1 H_2 \rangle = 0 \]
Values of $\mu^2$ in the landscape (quantum)

$V = \Lambda_{QCD}^3 h_2^0 + m_{H_2}^2 (h_2^0)^2 + \lambda (h_2^0)^4$

- $h_2^0 \sim \frac{\Lambda_{QCD}^3}{m_{H_2}^2}$
- $m_{H_2} \sim \Lambda_{QCD}$
- $h_2^0 \sim \frac{\Lambda_{QCD}}{\lambda^{1/3}}$

$\langle H_1 H_2 \rangle = 0$

$\Lambda_{QCD}^2$
\[ \mu^2 = \langle H_1 H_2 \rangle \]

\[ V = \Lambda_{QCD}^3 \hat{h}_2^0 + m_{H_1}^2 (h_2^0)^2 + \lambda (h_2^0)^4 \]

\[ m_{H_1}^2 < 0 \]

\[ \text{Triggering parameter} \]

\[ \text{quantum effects (QCD)} \]
\[ \mu^2 = \langle H_1 H_2 \rangle \]

\[ m_{H_1}^2 < 0 \]

\[ V = \Lambda_{QCD}^3 h_2^0 + m_{H_2}^2 (h_2^0)^2 + \lambda (h_2^0)^4 \]

\[ \mu_B^2 \]

\[ \mu_S^2 \]

\[ \Lambda_{QCD} \langle H_1 \rangle \]

\[ -\Lambda_{QCD} \]

\[ 0 \]

\[ \Lambda_{QCD} \]

\[ \text{Sign}(m_{H_2}^2) \sqrt{|m_{H_2}^2|} \]

Universe including ours

quantum effects (QCD)

classical
Strange universe with $\Lambda \sim (\frac{v}{\Lambda_H})^4 \Lambda_{us}$ for the atom formation.

$$\mu^2 = \langle H_1 H_2 \rangle$$

$$m^2_{H_1} < 0$$

$$V = \Lambda_{QCD}^3 h_2^0 + m^2_{H_2} (h_2^0)^2 + \lambda (h_2^0)^4$$

$\Lambda_{QCD} \langle H_1 \rangle$

$\mu^2_B$

$\mu^2_S$

$\mu^2 \approx \sqrt{|m^2_{H_1}| \frac{\Lambda_{QCD}^3}{m^2_{H_2}}}$

$\sqrt{|m^2_{H_2}|}$

$\Lambda_{QCD}$

$-\Lambda_{QCD}$

$0$ classical
Weak scale as a trigger

\[ \mu^2 \equiv \langle \mathcal{O} \rangle \]

as a function of \( m_h^2 \)

\[ \frac{\Lambda}{M_4^4} \frac{m_{H_1}^2}{\Lambda_H^2} \frac{m_{H_2}^2}{\Lambda_H^2} = \frac{1}{N_{\text{UV}}} \]

\[ \frac{M_4^4}{N_{\text{UV}}} \]

cosmological constant

small cosmological constant

\[ \mu^2 \ll \mu_S^2 \ll \mu_B^2 \]

weak scale
small $\mu^2$ is achieved by tuning in the landscape

$$\frac{\Lambda}{M_4^4} \frac{m_{H1}^2}{\Lambda_H^2} \frac{m_{H2}^2}{\Lambda_H^2} = \frac{1}{N_{UV}}$$

Cosmological constant

small cosmological constant

$$\Lambda = \frac{M_4^4}{N_{UV}} \frac{\Lambda_H^4}{\mu^4}$$

$\mu_S^2 \ll \mu^2 \ll \mu_B^2$

$\mu^2$ weak scale
Two big problems in theoretical physics

A. Cosmological constant problem
B. Gauge hierarchy problem

\[ \Lambda \sim 10^{-120} M_{Pl}^4 \]
\[ v^2 \sim 10^{-30} M_{Pl}^2 \]

We need at least \(10^{150}\) vacua to explain two small parameters.
We now couple to the energy landscape to provide a viable dark matter candidate. Due to the smaller annihilation cross section from weak interactions, the WIMP mass has a much higher energy density in universes with a larger Higgs vev (provided that they are reheated above the WIMP's weak scale). The low energy sector is generated by fields of mass close to the cutoff, and the number of non-degenerate minima depends on the Higgs vev. When the energy sector is generated by fields of mass close to the cutoff, and that the baryon energy density has come to dominate over the dark matter energy density, i.e. the baryon-to-dark-matter ratio is observed, our idea was discussed in Ref. [1].

Simple mechanisms that solve this problem by giving a baryon asymmetry inversely proportional to the Higgs vev can scan the 

\[ \Lambda_{\text{obs}} \]

\[ V_\Phi \]

\[ m_\Phi \sim M_* \]

\[ \langle \Phi \rangle \sim M_* \]

\[ V_\Phi = \sum_{i=1}^{N_1} \left[ \frac{M_i^2}{2} \Phi_i^2 + A_i \Phi_i^3 + \frac{\lambda_i}{4} \Phi_i^4 + M_{H,i}^2 \frac{\Phi_i}{M_*} |H|^2 + \ldots \right] \]

\[ m_H^2 \text{ scans} \]
$\mathcal{N}_{\text{UV}} \ll 10^{150}$

$v^2 \sim 10^{-30}$

$\Lambda_{cc} \sim 10^{-120}$

not possible in the HE landscape
Low Energy Landscape

\[ V^{(I)} = \sum_i \epsilon^2 (\phi_i^2 - M_*^2)^2 + \epsilon \kappa M_* \phi_i H_1 H_2 + V_H^{(I)}, \]

\[ V_H^{(I)} = (m_1^2)^{(I)} |H_1|^2 + (m_2^2)^{(I)} |H_2|^2 + \text{quartics} + q \lambda_u u^c H_1^* + q \lambda_d d^c H_1 + l \lambda_e e^c H_1 + \Lambda^{(I)} \]

\[ \langle h \rangle = 0 \quad \langle h \rangle \approx v \quad \langle h \rangle \gg v \]
Low energy landscape

A set of light scalar fields with degenerate vacua

$\phi O$ can break the degeneracy

Extra scanning of the cc is triggered by the weak scale

$$\mu^2 = \langle O \rangle$$

$\mu > \mu_B$ \quad \text{the other vacuum disappears}

$\mu < \mu_S$ \quad \text{hyperfine splitting of the scanning and it doesn’t help to reduce the cc}$
Values of the CC in the Landscape

\[ m_\phi \sim v^2 / M_* \quad \langle \phi \rangle \sim M_* \]

\[ \langle h \rangle = 0 \quad \langle h \rangle \sim v \quad \langle h \rangle \gg v \]

\[ V_\phi = \sum_{i=1}^{N_2} \frac{\epsilon_i^2}{4} \left( \phi_i^2 - M_*^{2,i} \right)^2 + \left( \sum_{i=1}^{N_2} \frac{\kappa_i \epsilon_i M_*^{3-\Delta_T}}{\sqrt{N_2}} \phi_i \mathcal{O}_T + \text{h.c.} \right) \]
the minima of Eq. (67) are lost at the scale to the vev transition,
smaller than $10^2$ meV.

We imagine that one of the usual mechanisms solves the strong CP problem, leaving at low energy a residual
the low energy landscape loses all its minima but one and from the point of view of the CC we have the same problem as in the vacua of the low energy landscape (right panel of Fig. 5).

Values of the Cosmological Constant in our two-sectors landscape. When $\langle h \rangle = 0$, $\langle h \rangle < \mu_S$, $\langle h \rangle \simeq \mu_S$, $\mu_S \lesssim \langle h \rangle \lesssim \mu_B$, $\langle h \rangle \gg \mu_B$.

High Energy Landscape

Low Energy Landscape
\[ \mathcal{O}_H = \kappa \epsilon M_* \phi H_1 H_2 \]
\[ V_{\text{loop}} \sim \kappa^2 \epsilon^2 M_*^2 \phi_i \phi_j \]
\[ \kappa^2 \epsilon^2 M_*^4 \ll \Lambda_* = \frac{M_*^4}{\mathcal{N}_{\text{UV}}} \]

To prevent IR scan from \( V_{\text{loop}} \)

\[ \Lambda(\mu^2) = \frac{\Lambda_H^4}{|m_{H_1}^2 m_{H_2}^2|} \Lambda_* = \frac{\Lambda_H^4}{\mu^4} \Lambda_* \quad \text{CC from UV scan} \]

Splitting should be larger than the CC from UV scan

\[ \frac{\Lambda_H^8 M_*^4}{v^{12}} \ll \mathcal{N}_{\text{UV}} \ll 10^{120} \frac{\Lambda_H^2}{v^2} \quad \rightarrow \quad \Lambda_H \ll 10^{12} \text{ GeV} \]

\[ \kappa \ll \frac{\mu_H^2}{\Lambda_H^2} \sim \frac{v^2}{\Lambda_H^2} \quad \rightarrow \quad \langle \mathcal{O}_H \rangle \sim \kappa v^4 \ll v^4 \]
Weak scale as a trigger

\[ \mu^2 \equiv \langle \mathcal{O} \rangle \]

as a function of \( m_h^2 \)

\[ \Lambda_{\text{min}} \propto \frac{1}{\mu^4} \]

\[ \frac{\Lambda}{M_*^4} \frac{m_{H_1}^2}{\Lambda_H^2} \frac{m_{H_2}^2}{\Lambda_H^2} = \frac{1}{N_{\text{UV}}} \]

\[ \frac{M_*^{4}}{N_{\text{UV}}} \]

\[ \mu^2 \ll \mu_S^2 \ll \mu_B^2 \]

small cosmological constant

cosmological constant

weak scale
A possibility of entirely different universe

\[ m_{H_1}^2 < 0 \quad m_{H_2}^2 > 0 \]

weak scale \[ \mu^2 = \Lambda_H \frac{\Lambda_{QCD}^3}{m_{H_2}^2} \]

cc \[ \Lambda(\mu^2) = \frac{\Lambda_H^2}{m_{H_2}^2} \frac{M^4_*}{N_{UV}} = \frac{\Lambda_H \mu^2}{\Lambda_{QCD}^3} \frac{M^4_*}{N_{UV}} \]

\[ \kappa^2 \mu_B^2 \mu^2 > \Lambda(\mu^2) \]

condition for IR scan \[ N_{UV} > \frac{1}{\kappa^2} \frac{\Lambda_H M^4_*}{\Lambda_{QCD}^3 \mu_B^2} \]

Fermion mass \[ \frac{\Lambda_{QCD}^3}{m_{H_2}^2} = \frac{\mu^2}{\Lambda_H} \leq \frac{v^2}{\Lambda_H} \quad \text{smaller at least by} \quad \frac{v}{\Lambda_H} \]

The cc should be smaller by \( \left( \frac{v}{\Lambda_H} \right)^4 \) for atoms to form
The weak scale as a trigger

I. Type 0 2HDM
FIG. 5. Experimental Constraints on the CP-even Higgs $H$ for $m_{H^\pm} = m_A$ and different values of $\Delta m$. From top to bottom we increase $m_{H^\pm}$. From left to right we move from 1% tuning ($\Delta m = 0.01 | \Delta m + 5 |$) to natural values of the quartics ($\Delta m = | \Delta m + 4 |$). In red we show the bound from $e^+e^- \to ZH$ at LEP \cite{3} and in yellow from $HZ$ associated production \cite{3} followed by decays to fermions. In light blue we display the current sensitivity of $H \to \ell^+\ell^-$ at LEP and the LHC \cite{23–25} and a projection for the HL-LHC obtained rescaling \cite{25}. In light green we show bounds from searches for $B \to K(\tau\nu)$ \cite{26,27} and $H \to K(\tau\nu)\mu\mu$ at LHCb \cite{26,27}. Indirect constraints from Higgs coupling measurements (purple and blue) are discussed in Section III B 2. The pink shaded area shows the strongest bound point-by-point between searches for flavor changing processes, mainly $b \to s$ \cite{7,8}, and LHC searches for $t \to H b$ \cite{9–12}. Theoretical constraints (in gray) from low energy Landau poles and the SM Higgs mass are summarized at the beginning of Section III B.
Domain wall from Z2 symmetry of H1

\[ \kappa \epsilon M_* \langle \phi \rangle H_1 H_2 \]

\[ B \mu_{\text{eff}} = \kappa \epsilon M_* \langle \phi \rangle \sim \kappa v^2 \]

Spontaneous breaking of Z2 from phi misalignment

Domain wall energy density starts to dominate at

\[ T \sim \left( \frac{v}{M_{Pl}} \right)^{1/2} v \sim \text{keV} \]

\[ \frac{B \mu_{\text{eff}} v^2}{v^3} \sim H \]

Biased potential annihilates domain walls

No domain wall problem for \( B \mu_{\text{eff}} \geq \frac{v^4}{M_{Pl}^2} \)
Scalar dark matter from the electroweak phase transition

Misalignment of the light scalar provides a dark matter

\[ \ddot{\phi} + 3H \dot{\phi} + V,\phi + c\langle \phi \rangle_T \langle H_1 H_2 \rangle_T = 0 \]

The last term provides a kick to the light scalar at EWPT

\[ \Delta \phi \sim \mathcal{O}(M_*) \]

The relic density is determined from EWPT
**Light Scalar Dark Matter**

![Graph](image)

\( m_\phi \geq 10^{-21} \text{eV} \)

Astrophysical constraints on fuzzy dark matter

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Planck's measurement of the most constraining scale leaves the horizon. Isocurvature perturbations from superradiance are cut off in the power spectrum of the Higgs boson in the cooling equivalence principle. The bound on the DM mass is still the subject of active research and goes beyond the scope of this work. Imposing the same requirement on the Jeans length induced by the dark matter means that Hubble is evaluated when the subscript is smaller than the typical size of a galaxy and is subsequently frozen. In quoting the bound we have used the subscript 0 and assumed isocurvature perturbations that set a mild constraint on Hubble during inflation. The bound \( \kappa \) on the trilinear coupling of a scalar coupled to the Higgs boson in the cooling equivalence principle goes beyond the scope of this work.

In addition to the laboratory and astrophysical constraints shown in the Figure, Planck's measurement of the most constraining scale leaves the horizon. Isocurvature perturbations from superradiance are cut off in the power spectrum of the Higgs boson in the cooling equivalence principle. The bound on the DM mass is still the subject of active research and goes beyond the scope of this work. Imposing the same requirement on the Jeans length induced by the dark matter means that Hubble is evaluated when the subscript is smaller than the typical size of a galaxy and is subsequently frozen. In quoting the bound we have used the subscript 0 and assumed isocurvature perturbations that set a mild constraint on Hubble during inflation. The bound \( \kappa \) on the trilinear coupling of a scalar coupled to the Higgs boson in the cooling equivalence principle goes beyond the scope of this work.
Summary

The smallness of the cc and the observed weak scale might have a tight connection in the landscape.

In the friendly landscape in which only the dimensionful parameters scan, the big landscape the cc scan might be sparse.

Electroweak symmetry breaking might break the degeneracy of light scalar vacua and can further scan the cc down to small one.

For the mechanism to work, (type 0) 2HDM is predicted and we would expect to discover additional Higgs bosons at the LHC.

Lots of light scalars can provide an excellent candidate of dark matter from their coherent oscillations (misalignment is made at the electroweak phase transition).
Backup
Basics of Type 0 2HDM

CP odd Higgs $A$ : PQ Goldstone boson

$$m_A^2 = -\lambda_5 v^2 \quad \lambda_5 < 0$$

CP even neutral Higgs $h$ and $H$ : $m_H \leq m_h$

$$\begin{pmatrix}
\lambda_1 v_1^2 & \lambda_{345} v_1 v_2 \\
\lambda_{345} v_1 v_2 & \lambda_2 v_2^2
\end{pmatrix}$$

$v_1 \ll v_2$

$$m_h^2 = \lambda_2 v_2^2 \quad \text{SM-like}$$

$$m_H^2 = (\lambda_1 - \frac{\lambda_{345}^2}{\lambda_2}) v_1^2 \quad \text{H lighter than } h$$

charged Higgs

$$m_{H^\pm}^2 = -\frac{\lambda_4 + \lambda_5}{2} v^2$$
Basics of Type 0 2HDM

$$\lambda_{345} = \lambda_3 + \lambda_4 + \lambda_5$$

$$g_{H\psi\psi} \approx -g_{h\psi\psi}^{\text{SM}} \frac{\lambda_{345}}{\lambda_2} \frac{v_1}{v}$$

Fermion couplings of H have double suppression

Fermio-phobic H

$$\lambda_{345} = 0$$

$$g_{H\psi\psi} = 0$$
Basics of Type 0
2HDM

\[ g_{H+tb^c} \simeq g_{htt}^{\text{SM}} \frac{v_1}{v} \]

\[ g_{H-tb^c} \simeq g_{hbb}^{\text{SM}} \frac{v_1}{v} \]

\[ g_{A\psi\psi} \simeq \pm g_{h\psi\psi}^{\text{SM}} \frac{v_1}{v} \]

\[ g_{HVVV} \simeq g_{hVVV}^{\text{SM}} \lambda_2 |1 - \frac{\lambda_{345}}{\lambda_2}| \frac{v_1}{v} \]

Gauge phobic H

\[ \lambda_{345} = \lambda_2 \]

\[ g_{HVVV} = 0 \]
Basics of Type 0
2HDM

\[ \lambda_{345} = \lambda_3 + \lambda_4 + \lambda_5 \]

\[ g_{HV} \sim g^{\text{SM}}_{hVV} \frac{|\lambda_2 - \lambda_{345}|}{\lambda_2} \frac{v_1}{v} \]

\[ \lambda_{hHH} \sim \lambda_{345}v \]

\[ \lambda_{hAA} \sim (\lambda_{345} - 2\lambda_5)v \]

\[ g_{AVVV} = 0 \]

\[ g_{ZAH} \sim -\frac{g}{2\cos\theta_W}(p_A + p_H) \] independent of \( \lambda_i \)
Basics of Type 0
2HDM

\[
\frac{g_{hVV} - g_{hVV}^{\text{SM}}}{g_{hVV}^{\text{SM}}} \approx -\frac{v_1^2}{2v^2} \left(1 - \frac{\lambda_{345}v^2}{m_h^2 - m_H^2}\right)^2
\]

\[
\frac{g_{h\psi\psi} - g_{h\psi\psi}^{\text{SM}}}{g_{h\psi\psi}^{\text{SM}}} \approx -\frac{v_1^2}{2v^2} \left(1 - \frac{\lambda_{345}^2v^4}{(m_h^2 - m_H^2)^2}\right)
\]

The deviation can be made to be small by choosing the ratio to be 1

\[
\lambda_{345}v^2 = m_h^2 - m_H^2
\]
Basics of Type 0
2HDM

CP even Higgs $H$ is predicted to be lighter than $h$

\[ \Lambda_{QCD}^2 \leq m_H^2 \leq m_h^2 \]

Charged Higgs and CP odd Higgs are lighter than 200 GeV

\[ m_{H^\pm}, m_A \leq 250 \text{ GeV} \]

\[ 175 \text{ GeV} \]

\[ \Lambda_{UV} = 500 \text{ GeV} \]
\[ \Lambda_{UV} = 10^7 \text{ GeV} \]

UV cutoff from Landau pole
The scale of Landau pole depending on the couplings

\[ \Lambda_{\text{UV}}(\text{GeV}) \]

\[ \Lambda_{\text{UV}}(\text{GeV}) = \lambda_1(\mu=m_\text{Z}) \]

Landau Pole

Stability
FIG. 1. Experimental Constraints on the type-0 2HDM in the $m_{H^\pm}$-$m_H$ plane (masses of the charged and new CP-even Higgses). The three panels correspond to three different choices for the vev of the new Higgs doublet $v_1 = h_{H^01} = (0.2, 0.3, 0.5)$. The mass of the new CP-Odd Higgs $m_A$ is fixed to 160 GeV. In all three cases we allow a 10% tuning of quartics, i.e. we take $|3| + |4| + |5| = 0.1$.

In red we show the bound from $H^\pm$ pair production at LEP [1, 2], and in yellow from $HZ$ associated production and decays to fermions [3]. In light blue we display the bound from ElectroWeak Precision Tests [4, 5] on the $S$, $T$ and $U$ oblique parameters [5, 6]. In light green we show bounds from searches for $B\to X_s$ [7, 8]. Indirect constraints from Higgs coupling measurements set an upper bound on BR($h\to HH$) (in blue). The impact of LHC searches for $t\to Hb$ is shown in pink [9–12]. Theoretical constraints (in gray) from low energy Landau poles and the SM Higgs mass are summarized at the beginning of Section III B.

At high masses there is no solution for the quartics which gives $m_h = 125$ GeV. But to our great surprise, it is not (yet) entirely excluded by collider searches! Some representative region of parameter space that is still viable, in the plane of new CP even and charged Higgs masses, is shown in Fig. 1, and the collider bounds will be discussed at greater length in Section III. There is a reasonable (if admittedly moderately tuned) region of parameter space where the new Higgs states have thus far escaped detection. This model will incisively be probed in the high-luminosity run of the LHC, and the new states can also be copiously produced at Higgs factories. If these new light states are seen, and the associated fingerprint of the $Z_4$ symmetry is confirmed, that would give direct experimental evidence for the "weak scale as a trigger".

Given some operator $O$ triggered by the weak scale, it is natural to try to use this trigger to attack the hierarchy problem in a new way. For instance, we can look for cosmological vacuum selection scenarios, that force $\mu^2 = h_{O1}^2$, (1) in the range $\mu^2_S \sim \mu^2 < \mu^2_B$, (2). Using $O_{H1} = H_1 H_2$ of our type-0 2HDM, this would force tuning for light Higgses. We present one such vacuum selection mechanism in the context of the landscape [1]. We give a field-theoretic model for the landscape, with a "UV landscape" containing moderately many vacua, not enough to find vacua with our small cosmological constant (CC). But we also imagine a separate "IR landscape", with $n$ ultra-light, weakly coupled scalars $i$, each with a (spontaneously broken) $Z_2$ discrete symmetry potentially giving a factor of 2 more vacua. The $i$ also couple to $O_{H1}$. If $h_{O1}$ is too small, the $2n$ vacua of the $i$ sector are all degenerate and they don’t help with making smaller vacuum energies possible. If $h_{O1}$ is too big, the symmetry is broken so badly that only one vacuum remains for each $i$, and there is again no way to find small vacuum energy. The only way to 1...
FIG. 5. Experimental Constraints on the CP-even Higgs $H$ for $m_{H^\pm} = m_A$ and different values of $\lambda_{345}$. From top to bottom we increase $m_{H^\pm}$. From left to right we move from 1% tuning ($\lambda_{345} = 0.01 | 4 + 5 |$) to natural values of the quartics ($\lambda_{345} = | 4 + 5 |$). In red we show the bound from $e^+e^-\rightarrow ZH$ at LEP [3] and in yellow from $HZ$ associated production [3] followed by decays to fermions. In light blue we display the current sensitivity of $H \rightarrow \mu\mu$ at LEP and the LHC [23–25] and a projection for the HL-LHC obtained rescaling [25]. In light green we show bounds from searches for $B \rightarrow K(\tau\nu)$ at LHCb [26, 27]. Indirect constraints from Higgs coupling measurements (purple and blue) are discussed in Section III B 2. The pink shaded area shows the strongest bound point-by-point between searches for flavor changing processes, mainly $b \rightarrow s$ [7, 8], and LHC searches for $t \rightarrow H + b$ [9–12]. Theoretical constraints (in gray) from low energy Landau poles and the SM Higgs mass are summarized at the beginning of Section III B.

To a few percent. The bound is shown in pink in Fig.s 1 and 5. When the mass difference between $H^\pm$ and $A$ becomes $O(\pm W)$, also the decays $H^\pm \rightarrow W^\pm A$ become relevant [29], but we do not consider this parameter space in our analysis since it is disfavored by bounds on $h \rightarrow AA$ and Electroweak precision measurements.

For larger $m_{H^\pm}$, when top decays are not kinematically allowed, direct searches for $H^\pm$, that typically target $pp \rightarrow \bar{t}bH +$ and decays to $tb$ and $\nu\nu$, are not yet sensitive to our parameter space [9, 30–37]. To conclude this brief overview of the charged Higgs, it is interesting to notice that a CMS search for stau pair production [38] has a sensitivity comparable to LHC searches targeting $H^+ \rightarrow$ single production. The latter can be decoupled.
Scalar dark matter from the electroweak phase transition

Misalignment of the light scalar provides a dark matter

\[ \ddot{\phi} + 3H \dot{\phi} + V,\phi + c\langle \phi \rangle_T \langle H_1 H_2 \rangle_T = 0 \]

Misalignment of the light scalar provides a dark matter

1. There is a misalignment of the light scalar after inflation

2. When \( H \sim m_\phi \), the scalar starts the oscillation

3. **Electroweak phase transition** also gives a misalignment
Scalar dark matter from the electroweak phase transition

Misalignment of the light scalar provides a dark matter

$$\ddot{\phi} + 3H\dot{\phi} + V,\phi + c\langle \phi \rangle_T \langle H_1 H_2 \rangle_T = 0$$

The last term provides a kick to the light scalar at EWPT

$$\Delta \phi \sim \mathcal{O}(M_*)$$

The relic density is determined from EWPT
A kick to the light scalar at EWPT

\[
\kappa^2 \mu_B^4 \frac{\Delta \phi}{M_*} \sim \kappa^2 \mu^2 \mu_B^2
\]

\[
\frac{\Delta \phi}{M_*} \sim \frac{\mu^2}{\mu_B^2}
\]

When we are close to the upper bound on \( \mu \),

\[
\mu \sim \mu_B \\
\Delta \phi \sim M_* \\
\kappa^2 \mu_B^4 \sim \kappa^2 v^4
\]
$m_\phi > H(v)$
\( m_\phi < H(v) \)
Sketch for the relic abundance of light scalar dark matter

At the EWPT, the amount of the misaligned energy density: \( \kappa^2 v^4 \)

current dark matter density: \( \frac{v^8}{M_{Pl}^4} \)

EWPT to matter radiation equality: \( \frac{T_{eq}}{T_W} \sim \frac{v^2}{M_{Pl}} = \frac{v}{M_{Pl}} \)

scalar oscillation: \( \frac{1}{a^3} \)
radiation: \( \frac{1}{a^4} \)
Sketch for the relic abundance of light scalar dark matter

\[ m_\phi > H(v) \quad \Rightarrow \quad 10^{-5} \text{ eV} \]

\[ \kappa \sim \sqrt{\frac{v}{M_{Pl}}} \sim 10^{-8} \]

\[ m_\phi < H(v) \]

\[ \kappa^2 v^4 \left( \frac{T_{eq}}{T_{osc}} \right)^3 \sim \frac{v^8}{M_{Pl}^4} \]

\[ T_{eq} \sim \frac{v^2}{M_{Pl}} \quad T_{osc} \sim \sqrt{m_\phi M_{Pl}} \]

\[ \kappa \sim \left( \frac{M_*}{M_{Pl}} \right)^{1/4} \sqrt{\frac{v}{M_{Pl}}} \lesssim 10^{-8} \]
are completely uncorrelated with curvature perturbations. The most constraining scale measured by Planck is when the perturbation leaves the horizon at scales of $O(10^{15} \text{Mpc})$ at all times between today and matter radiation equality. This is satisfied in all our DM parameter sets. However, since we neglect unimportant terms in the potential, these anharmonic terms in the potential do not lead to a suppression of the power of isocurvature perturbations.}

We have obtained our previous results neglecting the anharmonic terms in the potential. These anharmonic terms in the potential do not lead to a suppression of the power. The subscript $i$ sets a mild constraint on Hubble during inflation that at fixed $m_{\phi} = 10^{-15} \text{eV}$, this is consistent with observational bounds on the lightest viable DM mass.

In Fig. 9, we also show laboratory and astrophysical constraints on isocurvature perturbations. In addition to the laboratory and astrophysical constraints shown in the Figure, Planck's measurement sets a mild constraint on Hubble during inflation that at fixed $m_{\phi} = 10^{-15} \text{eV}$, this is consistent with observational bounds on the lightest viable DM mass.

In Fig. 9, we also show laboratory and astrophysical constraints on isocurvature perturbations. In addition to the laboratory and astrophysical constraints shown in the Figure, Planck's measurement sets a mild constraint on Hubble during inflation that at fixed $m_{\phi} = 10^{-15} \text{eV}$, this is consistent with observational bounds on the lightest viable DM mass.

Astrophysical constraints on light scalar dark matter. They include tests of the equivalence principle and black hole superradiance. The light pink shaded area at the bottom of the plot is the theoretical constraint on cooling. The bound on the DM mass is still the subject of active research and goes beyond the scope of this work. The subscript $i$ sets a mild constraint on Hubble during inflation that at fixed $m_{\phi} = 10^{-15} \text{eV}$, this is consistent with observational bounds on the lightest viable DM mass.