Conformal versus Non-Conformal 2HDM: Phase Transitions and Gravitational Waves

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Work in progress

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Outline of the talk

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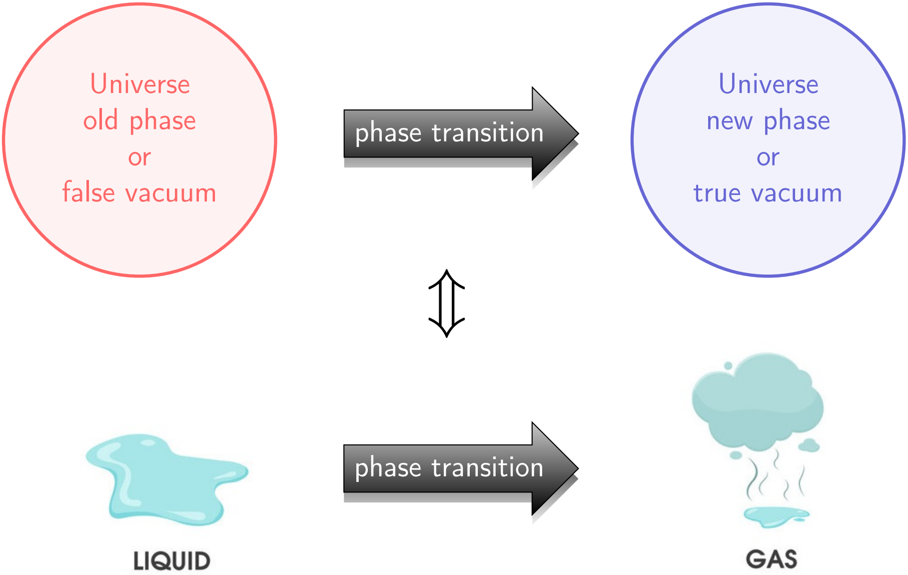
1. Introduction

2. Model

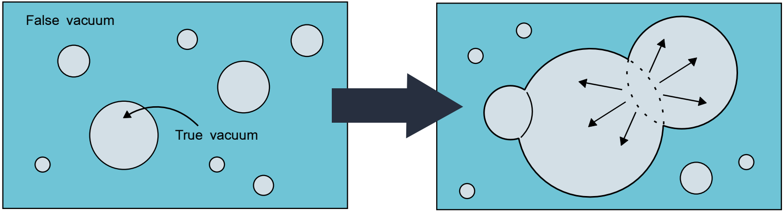
3. Results

4. Conclusion

Phase transition in the early Universe

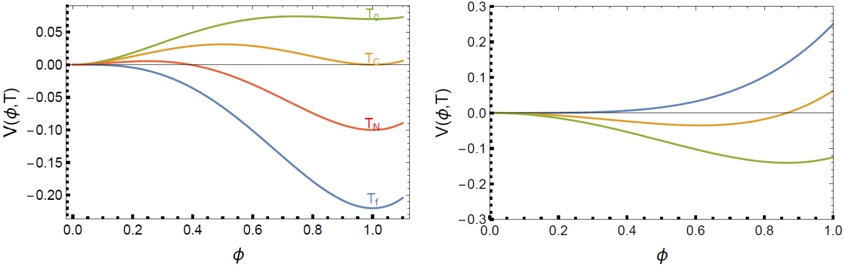


Bubbles of new phase and collision



$\Updownarrow$

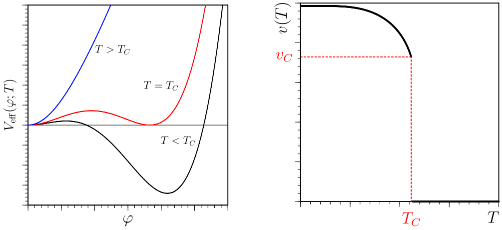
Order of phase transitions

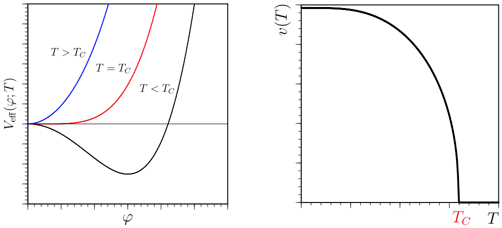


Mazumdar, White, arXiv:1811.01948

- 1^{st}-order phase transitions are abrupt (existence of a barrier): the order parameter (VEV ϕ) changes *discontinuously* from zero to a non-zero value
- 2^{nd}-order phase transitions/crossovers are smooth (no barrier): the order parameter (VEV ϕ) changes *continuously* from zero to a non-zero value

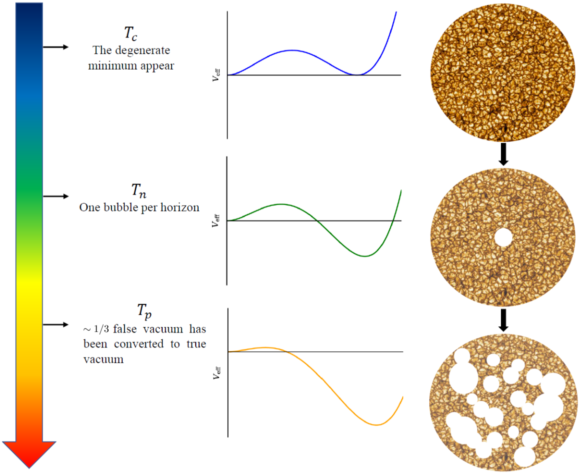
Order of phase transitions





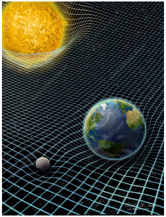
Senaha, Symmetry 2020, 12(5), 733

First-order phase transitions

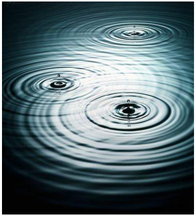


Wang, Huang, Zhang, arXiv:2003.08892

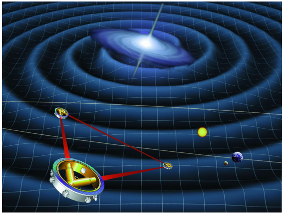
Gravitational waves



spacetime

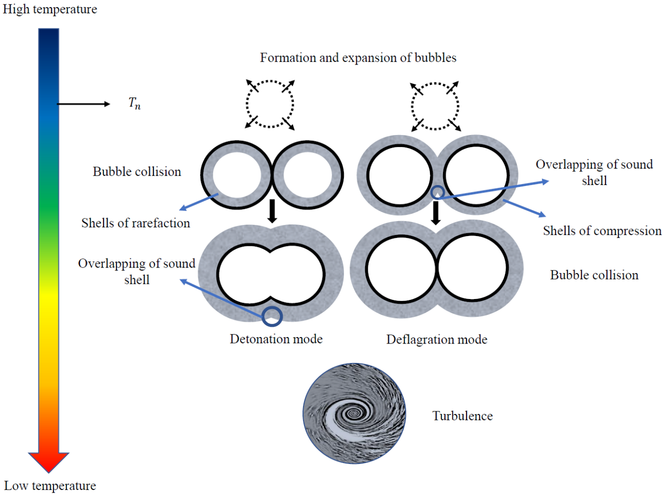


spacetime perturbation



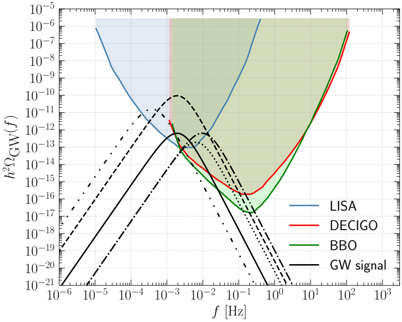
detection of gravitational waves

Gravitational-wave production



Wang, Huang, Zhang, arXiv:2003.08892

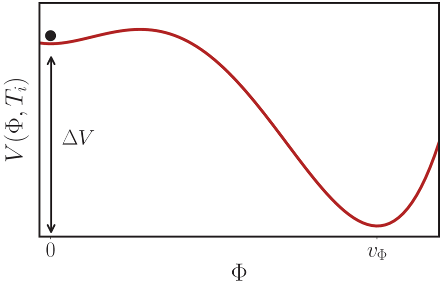
Gravitational-wave power spectrum



PLI sensitivity curves $h^2\Omega_{\text{PLI}}(f)$ for LISA, DECIGO and BBO, with $t_{\text{obs}} = 4$ yr and $\rho_{\text{thrs}} = 10$.

- - - - - : α increases
- · · · · · : β/H_* increases
- −·−·− : T_* increases, while keeping β/H_* constant
- −· ·−· ·− : T_* increases, considering $H_* \sim T_*^2$

Supercooled first-order phase transitions



Sagunski, Schicho, Schmitt, arXiv:2303.02450

For supercooled first-order phase transitions

- the Universes remains trapped in the false vacuum until $T \ll T_c$
- ΔV increases as T decreases
- for $T < T_i$, the Universe becomes vacuum dominated ($\Delta V = \frac{\pi^2}{30} g_* T_i^4$) and enters a period of thermal inflation
- percolation may never happen if thermal inflation is too important $\rightarrow$ the phase transition may never complete

Outline of the talk

Motivations

- 2HDM is motivated by MSSM
- nearly conformal dynamics typically lead to large supercooling
- compare the results from phase transitions and gravitational waves between the 2HDM and its conformal version

C2HDM tree-level potential

The tree-level potential of the classically conformal CP-conserving 2HDM is given by

$$V_0 = \frac{\lambda_1}{2}(\Phi_1^\dagger\Phi_1)^2 + \frac{\lambda_2}{2}(\Phi_2^\dagger\Phi_2)^2 + \lambda_3\Phi_1^\dagger\Phi_1\Phi_2^\dagger\Phi_2 + \lambda_4\Phi_1^\dagger\Phi_2\Phi_2^\dagger\Phi_1 + \frac{\lambda_5}{2}\left[(\Phi_1^\dagger\Phi_2)^2 + (\Phi_2^\dagger\Phi_1)^2\right],$$

where Φ_1 and Φ_2 denote the two SU(2)$_L$ Higgs doublets and where we have imposed $\mathbb{Z}_2$ discrete symmetries to avoid FCNC at tree-level.

C2HDM potential

The two SU(2)$_L$ Higgs doublets are given by

$$\Phi_1 = \frac{1}{\sqrt{2}} \begin{pmatrix} \rho_1 + i\eta_1 \\ v_1 + \phi_1 + i\psi_1 \end{pmatrix} \quad \text{and} \quad \Phi_2 = \frac{1}{\sqrt{2}} \begin{pmatrix} \rho_2 + i\eta_2 \\ v_2 + \phi_2 + i\psi_2 \end{pmatrix},$$

with $v_1 = v\cos\beta, v_2 = \sin\beta$ and $v \simeq 246$ GeV the VEV of the SM Higgs doublets

Tadpole conditions

Extremising V_0 along ϕ_1 and ϕ_2 direction leads to the following conditions:

$$\frac{\lambda_1}{\lambda_2} = \tan^4\beta, \qquad \sqrt{\lambda_1\lambda_2} = -\lambda_{345},$$

where $\lambda_{345} \equiv \lambda_3 + \lambda_4 + \lambda_5$.

Mass matrices

In the vacuum and using the tadpole conditions, the mass matrices, after diagonalisation, are given by

$$M_S^2 = \begin{pmatrix} m_H^2 & 0 \\ 0 & 0 \end{pmatrix}, \quad M_P^2 = \begin{pmatrix} m_A^2 & 0 \\ 0 & 0 \end{pmatrix}, \quad M_C^2 = \begin{pmatrix} m_{H^\pm}^2 & 0 \\ 0 & 0 \end{pmatrix},$$

with

$$m_H^2 = -\lambda_{345} v^2 = \sqrt{\lambda_1 \lambda_2} v^2, \quad m_A^2 = -\lambda_5 v^2, \quad m_{H^\pm}^2 = -\frac{1}{2}(\lambda_4 + \lambda_5) v^2.$$

Finding $m_h^2 = 0$ means the direction h is the flat direction (second derivative is zero)

Flat direction

The C2HDM is naturally aligned at tree-level: $\alpha = \beta - \pi/2$ such that M_S is diagonalised.

Flat direction along

$$\phi \equiv h = -\sin\alpha\phi_1 + \cos\alpha\phi_2 = \cos\beta\phi_1 + \sin\beta\phi_2$$

with

$$\phi_1 = \phi\cos\beta, \quad \phi_2 = \phi\sin\beta,$$
$$\langle\phi\rangle = v, \quad \langle\phi_1\rangle = v\cos\beta = v_1, \quad \langle\phi_2\rangle = v\sin\beta = v_2$$

Parametrisation

Using tadpole conditions and the relation of the mass matrices with the physical masses, we obtain the following parametrisation:

$$\lambda_1 = \frac{m_H^2}{v^2}\tan^2\beta, \quad \lambda_2 = \frac{m_H^2}{v^2\tan^2\beta}, \quad \lambda_3 = \frac{2m_{H^\pm}^2 - m_H^2}{v^2},$$
$$\lambda_4 = \frac{-2m_{H^\pm}^2 + m_A^2}{v^2}, \quad \lambda_5 = \frac{-m_A^2}{v^2},$$

for which $V_0(\phi) = 0$.

$\Longrightarrow$ along the flat direction, the scalar potential is generated at one-loop order .

Coleman-Weinberg potential

In Landau gauge, the Coleman-Weinberg potential V_{CW} is defined in the $\overline{\text{MS}}$ scheme as

$$V_{\mathsf{CW}}(\phi) = \frac{1}{64\pi^2} \sum_i (-1)^F g_i \, m_i^4(\phi) \left[\ln \frac{m_i^2(\phi)}{\mu^2} - C_i \right],$$

where

- $i \in \{t, W^\pm, Z, \gamma, h, H, G_0, A, G^\pm, H^\pm\}$
- g_i: the number of degrees of freedom
- $m_i^2(\phi)$: i^{th} eigenvalue of the field-dependent mass matrix $(m^2)_{ab} \equiv \partial^2 V / \partial\phi_a \partial\phi_b$
- μ: renormalisation scale
- $F = 1$ for fermions and $F = 0$ for bosons
- $C_i = 3/2$ for scalars, fermions and $C_i = 5/6$ for gauge bosons

One-loop improved tadpole conditions

Tree-level tadpole conditions:

$$T_{\phi_1} \equiv \left\langle \frac{\partial V_0}{\partial \phi_1} \right\rangle = v_1 \left(\lambda_1 v_1^2 + \frac{1}{2} \lambda_{345} v_2^2 \right) = 0$$

$$T_{\phi_2} \equiv \left\langle \frac{\partial V_0}{\partial \phi_2} \right\rangle = v_2 \left(\lambda_2 v_2^2 + \frac{1}{2} \lambda_{345} v_1^2 \right) = 0$$

$\Rightarrow$ one-loop improved tadpole conditions:

Lee, Pilaftsis, arXiv:1201.4891

$$\left\langle \frac{\partial (V_0 + V_{\text{CW}})}{\partial \phi_i} \right\rangle = T_{\phi_i} + \frac{v_i v^2}{64\pi^2} \Delta \hat{t}_i = 0, \quad i = 1, 2$$

with

$$\begin{aligned} \Delta \hat{t}_i = & \frac{1}{v^2} \left[4\lambda_{345} m_H^2 \left(1 - \log \frac{m_H^2}{\mu^2} \right) + 4\lambda_5 m_A^2 \left(1 - \log \frac{m_A^2}{\mu^2} \right) + 4\lambda_{45} m_{H\pm}^2 \left(1 - \log \frac{m_{H\pm}^2}{\mu^2} \right) \right. \\ & \left. - 6 g_2^2 m_W^2 \left(\frac{1}{3} - \log \frac{m_W^2}{\mu^2} \right) - 3(g_1^2 + g_2^2) m_Z^2 \left(\frac{1}{3} - \log \frac{m_Z^2}{\mu^2} \right) + 24 y_f^2 m_t^2 \left(1 - \log \frac{m_t^2}{\mu^2} \right) \delta_{Ii} \right] \end{aligned}$$

One-loop improved mass matrices

$$M_S^2 = \begin{pmatrix} \lambda_1 v_1^2 + \frac{T_{\phi_1}}{v_1} + \left\langle \frac{\partial^2 V_{\mathsf{CW}}}{\partial \phi_1^2} \right\rangle & \lambda_{345} v_1 v_2 + \left\langle \frac{\partial^2 V_{\mathsf{CW}}}{\partial \phi_1 \partial \phi_2} \right\rangle \\ \lambda_{345} v_1 v_2 + \left\langle \frac{\partial^2 V_{\mathsf{CW}}}{\partial \phi_1 \partial \phi_2} \right\rangle & \lambda_2 v_2^2 + \frac{T_{\phi_2}}{v_2} + \left\langle \frac{\partial^2 V_{\mathsf{CW}}}{\partial \phi_2^2} \right\rangle \end{pmatrix}$$

with

$$\left\langle \frac{\partial^2 V_{\mathsf{CW}}}{\partial \phi_i \partial \phi_j} \right\rangle = \frac{1}{64\pi^2}(v^2 \Delta \hat{t}_i \delta_{ij} + v_i v_j \Delta \hat{m}_{ij}^2),$$

$$\Delta \hat{m}_{ij}^2 \equiv 8\lambda_{345}^2 \log \frac{|m_H^2|}{\mu^2} + 8\lambda_5^2 \log \frac{m_A^2}{\mu^2} + 4\lambda_{45}^2 \log \frac{m_{H^\pm}^2}{\mu^2} + g_2^4 \left(2 + 3 \log \frac{m_W^2}{\mu^2}\right) + \frac{g_1^4 + g_2^4}{2} \left(2 + 3 \log \frac{m_Z^2}{\mu^2}\right) - 24 y_f^4 \log \frac{m_t^2}{\mu^2} \delta_{ij} \delta_{Ii}$$

One-loop improved mass matrices

$$M_S^2 = \begin{pmatrix} \lambda_1 v_1^2 + \frac{T_{\phi_1}}{v_1} + \frac{1}{64\pi^2}(v^2\Delta\hat{t}_1 + v_1^2\Delta\hat{m}_{11}^2) & \lambda_{345}v_1v_2 + \frac{1}{64\pi^2}v_1v_2\Delta\hat{m}_{12}^2 \\ \lambda_{345}v_1v_2 + \frac{1}{64\pi^2}v_1v_2\Delta\hat{m}_{12}^2 & \lambda_2 v_2^2 + \frac{T_{\phi_2}}{v_2} + \frac{1}{64\pi^2}(v^2\Delta\hat{t}_2 + v_2^2\Delta\hat{m}_{22}^2) \end{pmatrix}$$

- one-loop improved tadpole conditions: $\frac{T_{\phi_i}}{v_i} + \frac{1}{64\pi^2}v^2\Delta\hat{t}_i = 0$
- $\Delta\hat{m}_{ij}^2 = 0$ for M_P^2 and $M_C^2 \Longrightarrow (M_{P,C}^2)_{\text{one-loop}} = (M_{P,C}^2)_{\text{tree-level}}$
- M_S^2 has now a second non-zero eigenvalue $m_h^2 \neq 0$ as a consequence of the breaking of scale invariance at loop order

Coleman-Weinberg potential

Along the flat direction we have $m_i^4(\phi) = m_i^4(v)\phi^4/v^4$. We can then rewrite V_{CW} as

$$V_{\mathsf{CW}}(\phi) = \phi^4 \left(A + B \ln \frac{\phi^2}{\mu^2} \right),$$

with

$$A = \frac{1}{64\pi^2 v^4} \sum_i (-1)^F g_i\, m_i^4(v) \left[\ln \frac{m_i^2(v)}{\mu^2} - C_i \right],$$

$$B = \frac{1}{64\pi^2 v^4} \sum_i (-1)^F g_i\, m_i^4(v).$$

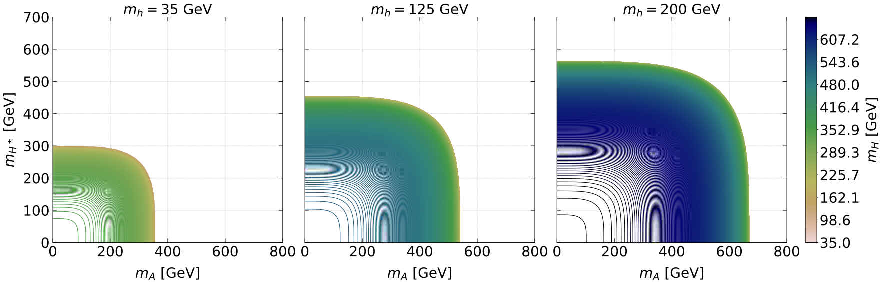
Gilderner-Weinberg renormalisation scale Λ_{GW}

By minimising V_{CW}, we find that $\mu = \Lambda_{\mathsf{GW}}$ is defined as

$$\ln \frac{\Lambda_{\mathsf{GW}}}{v} = \frac{A}{2B} + \frac{1}{4}.$$

At this scale Λ_{GW}, the minimum of V_{CW} is at $\phi = v$.

Bounds on heavy scalar masses



- sum rule: $m_H^4 = 8\pi^2 v^2 m_h^2 + 12 m_t^4 - 6 m_{W^\pm}^4 - 3 m_Z^4 - m_A^4 - 2 m_{H^\pm}^4$
- we need $m_H^4 > 0$ in order to have $m_H^2 \in \mathbb{R} \rightarrow$ upper bounds for m_H, m_A and $m_{H^\pm}$
- with the constraint $m_H > m_h \simeq 125$ GeV (otherwise perturbation theory breaks down), the sum rule is a reliable method to compute m_H

  Eichten, Lane, arXiv:2209.06632

One-loop thermal effective potential

We study the phase transitions in the early Universe: large temperature $\Rightarrow$ finite-temperature quantum field theory must be used to take thermal effects into account.

The resulting one-loop thermal effective potential in the C2HDM reads

$$V_{\text{eff}} = V_0 + V_{\text{CW}} + V_{\text{1L}}^T$$

- V_0: tree-level potential
- V_{CW}: Coleman-Weinberg potential
- V_{1L}^T: one-loop thermal corrections

One-loop thermal effective potential

$$V_{\mathsf{eff}} = V_0 + V_{\mathsf{CW}} + V^T_{\mathsf{1L}}.$$

The last term, V^T_{1L}, is defined as

$$V^T_{\mathsf{1L}}(\phi, T) = \frac{T^4}{2\pi^2} \sum_i g_i \, J_{\mathsf{B/F}} \left(\frac{m_i^2(\phi)}{T^2} \right) - \frac{\pi^2}{90} g'_* T^4$$

with

$$J_{\mathsf{B/F}}(y^2) = (-1)^F \int_0^\infty x^2 \log \left[1 \mp e^{-\sqrt{x^2+y^2}} \right]$$

- T: temperature
- g_i: degrees of freedom of the i^{th} particle
- $m_i^2(\phi)$: i^{th} eigenvalue of the field-dependent mass matrix $(m^2)_{ab} \equiv \partial^2 V / \partial\phi_a \partial\phi_b$
- g'_*: relativistic degrees of freedom of the remaining (light) particles

High-temperature expansion

In the high-temperature limit ($|y^2| = |m_i^2/T^2| \ll 1$), we can expand the thermal function J as following:

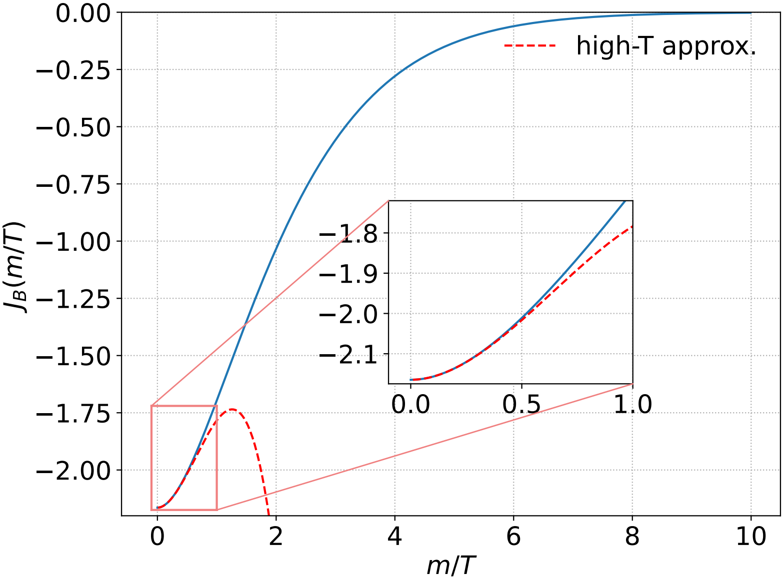
Curtin, Meade, Ramani, arXiv:1612.00466

$$J_B(y^2) \approx -\frac{\pi^4}{45} + \frac{\pi^2}{12}y^2 - \frac{\pi}{6}y^3 - \frac{1}{32}y^4 \log\left(\frac{y^2}{a_b}\right)$$
$$J_F(y^2) \approx -\frac{7\pi^4}{360} + \frac{\pi^2}{24}y^2 + \frac{1}{32}y^4 \log\left(\frac{y^2}{a_f}\right)$$

- $a_b = \pi^2 \exp(3/2 - 2\gamma_E)$
- $a_f = 16\pi^2 \exp(3/2 - 2\gamma_E)$

with $\gamma_E \simeq 0.577$, the Euler–Mascheroni constant.

High-temperature expansion



Imaginary contributions

The cubic term in

$$J_B(y^2) \approx -\frac{\pi^4}{45} + \frac{\pi^2}{12} y^2 - \frac{\pi}{6} y^3 - \frac{1}{32} y^4 \log\left(\frac{y^2}{a_b}\right)$$

can be imaginary for $m_i^2 < 0$. Indeed,

$$y^3 = (y^2)^{3/2} = \left(\frac{m_i^2}{T^2}\right)^{3/2}$$

$$\Rightarrow \operatorname{Im} y^3 \neq 0 \ \text{ for } \ m_i^2 < 0$$

Likewise for the $\ln m_i^2$ term in V_{CW}.

$\Rightarrow$ Thus we always consider the real part of V_{eff} in our calculations.

Parwani thermal resummation

Finally, we use the thermally improved finite-temperature potential, which is obtained by adding to the field-dependent masses in V_{CW} and V_{1L}^T the leading thermal corrections:

$$m_i^2(\phi) \to m_i^2(\phi) + c_i T^2,$$

where the coefficients c_i are given by

$$c_h = \frac{1}{16}(g_1^2 + 3g_2^2) + \frac{3\lambda_1 + 2\lambda_3 + \lambda_4}{12},$$
$$c_H = \frac{1}{16}(g_1^2 + 3g_2^2) + \frac{1}{4}y_t^2 + \frac{3\lambda_2 + 2\lambda_3 + \lambda_4}{12}$$

for the scalars.

NC2HDM tree-level potential

NC2HDM tree-level potential:

$$V_0 = m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - m_{12}^2 \left[\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1 \right] + \frac{\lambda_1}{2} (\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2} (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 \Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 + \lambda_4 \Phi_1^\dagger \Phi_2 \Phi_2^\dagger \Phi_1 + \frac{\lambda_5}{2} \left[(\Phi_1^\dagger \Phi_2)^2 + (\Phi_2^\dagger \Phi_1)^2 \right].$$

Parametrisation

In the alignment limit we have

$$\begin{aligned}
m_{11}^2 &= -\frac{m_h^2}{2} + m_{12}^2 \tan\beta, \\
m_{22}^2 &= -\frac{m_h^2}{2} + \frac{m_{12}^2}{\tan\beta}, \\
\lambda_1 v^2 &= m_h^2 + m_H^2 \tan^2\beta - M^2 \tan^2\beta, \\
\lambda_2 v^2 &= m_h^2 + \frac{m_H^2}{\tan^2\beta} - \frac{M^2}{\tan^2\beta}, \\
\lambda_3 v^2 &= m_h^2 - m_H^2 + 2m_{H^\pm}^2 - M^2, \\
\lambda_4 v^2 &= M^2 + m_A^2 - 2m_{H^\pm}^2, \\
\lambda_5 v^2 &= M^2 - m_A^2,
\end{aligned}$$

with

$$M^2 \equiv \frac{m_{12}^2}{\sin\beta\cos\beta}.$$

One-loop thermal effective potential

The one-loop thermal effective potential in the NC2HDM reads

$$V_{\mathsf{eff}} = V_0 + V_{\mathsf{CW}} + V_{\mathsf{CT}} + V_{\mathsf{1L}}^T$$

- V_0: tree-level potential
- V_{CW}: Coleman-Weinberg potential
- V_{CT}: counterterms
- V_{1L}^T: one-loop thermal corrections

One-loop thermal effective potential

For the treatment of the phase transitions, we suppose that excursions in the field space occur only along the direction ϕ ($\phi_1 = \phi \cos\beta$, $\phi_2 = \phi \sin\beta$).
$\Rightarrow$ the tree-level potential

$$
\begin{aligned}
V_0 = {} & m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - m_{12}^2 \left[\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1 \right] \\
& + \frac{\lambda_1}{2} (\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2} (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 \Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 \\
& + \lambda_4 \Phi_1^\dagger \Phi_2 \Phi_2^\dagger \Phi_1 + \frac{\lambda_5}{2} \left[(\Phi_1^\dagger \Phi_2)^2 + (\Phi_2^\dagger \Phi_1)^2 \right],
\end{aligned}
$$

becomes

$$
\begin{aligned}
V_0(\phi) = {} & \frac{1}{2} m_{11}^2 \phi^2 \cos^2\beta + \frac{1}{2} m_{22}^2 \phi^2 \sin^2\beta - m_{12}^2 \phi^2 \cos\beta \sin\beta \\
& + \frac{1}{8} \lambda_1 \phi^4 \cos^4\beta + \frac{1}{8} \lambda_2 \phi^4 \sin^4\beta + \frac{1}{4} \lambda_{345} \phi^4 \cos^2\beta \sin^2\beta .
\end{aligned}
$$

One-loop thermal effective potential

$$V_{\mathsf{eff}} = V_0 + V_{\mathsf{CW}} + V_{\mathsf{CT}} + V^T_{\mathsf{1L}}$$

The third term, V_{CT}, contains the finite parts of the counterterms that are fixed such that the scalar VEVs and masses remain at their tree-level values at the $T = 0$ global minimum (v_1, v_2):

$$V_{\mathsf{CT}}(\phi_1, \phi_2) = \delta m^2_{11}\phi^2_1 + \delta m^2_{22}\phi^2_2 + \delta\lambda_1\phi^4_1 + \delta\lambda_2\phi^4_2 + \delta\lambda_{345}\phi^2_1\phi^2_2,$$

such that the following renormalisation conditions are satisfied:

$$\frac{\partial V_{\mathsf{CT}}}{\partial \phi_i}\bigg|_{\mathsf{vev}} = -\frac{\partial V_{\mathsf{CW}}}{\partial \phi_i}\bigg|_{\mathsf{vev}}, \quad i = 1, 2,$$

$$\frac{\partial^2 V_{\mathsf{CT}}}{\partial \phi^2_i}\bigg|_{\mathsf{vev}} = \left(-\frac{\partial^2 V_{\mathsf{CW}}|_{G\equiv 0}}{\partial \phi^2_i} + \frac{1}{32\pi^2}\sum_{G=G^0,G^\pm}\left(\frac{\partial m^2_G}{\partial \phi_i}\right)^2 \ln\left(\frac{m^2_{\mathsf{IR}}}{\mu^2}\right)\right)\Bigg|_{\mathsf{vev}}, \quad i = 1, 2,$$

$$\frac{\partial^2 V_{\mathsf{CT}}}{\partial \phi_1 \partial \phi_2}\bigg|_{\mathsf{vev}} = \left(-\frac{\partial^2 V_{\mathsf{CW}}|_{G\equiv 0}}{\partial \phi_1 \partial \phi_2} + \frac{1}{32\pi^2}\sum_{G=G^0,G^\pm}\frac{\partial m^2_G}{\partial \phi_1}\frac{\partial m^2_G}{\partial \phi_2} \ln\left(\frac{m^2_{\mathsf{IR}}}{\mu^2}\right)\right)\Bigg|_{\mathsf{vev}}.$$

Theoretical constraint

Constraints on λ's from

- perturbativity
- perturbative unitarity
- bounded-from-below potential

Experimental constraints

Constraints on heavy scalar masses from

Arcadi, N.B., Djouadi, Kannike, arXiv:2212.14788

- LEP2 search: $m_A \gtrsim 90$ GeV ($e^+e^- \to hA$) and $m_{H^\pm} \gtrsim 80$ GeV ($e^+e^- \to H^+H^-$)
- constraint from FCNC via the loop-induced decay process $B \to X_s\gamma$ ($b \to s\gamma$):
  - type II: $m_{H^\pm} \gtrsim 800$ GeV for any $\tan\beta$
  - type I: $m_{H^\pm} \gtrsim 500$ GeV for $\tan\beta \gtrsim 1$
- EW precision measurements: mass splitting between m_H, m_A and $m_{H^\pm}$ cannot be too large

Scan of the parameter space

To analyse the phase transition dynamics in the C2HDM, we scan the parameter space over the following range:

$$m_A \in [90, 1000] \text{ GeV}, \quad m_{H^\pm} \in [80, 1000] \text{ GeV}, \quad \tan\beta \in [0.1, 50].$$

In the NC2HDM, in addition, we scan over

$$m_H \in [m_h, 1000] \text{ GeV}, \quad |m_{12}| \in [100, 1000] \text{ GeV}.$$

We then analyse the points which pass the aforementioned theoretical and experimental constraints.

Outline of the talk

PT parameters: definition

- inverse time duration β of PT normalised to the Hubble parameter:

$$\frac{\beta}{H_*} = (8\pi)^{1/3} \frac{v_w}{R_* H_*},$$

  with R_* the mean bubble separation
- PT strength α:

$$\alpha = \frac{1}{\frac{\pi^2}{30} g_* T_p^4} \left(\Delta V - \frac{T_p}{4} \Delta \frac{\partial V}{\partial T} \right) \Bigg|_{T_p}$$

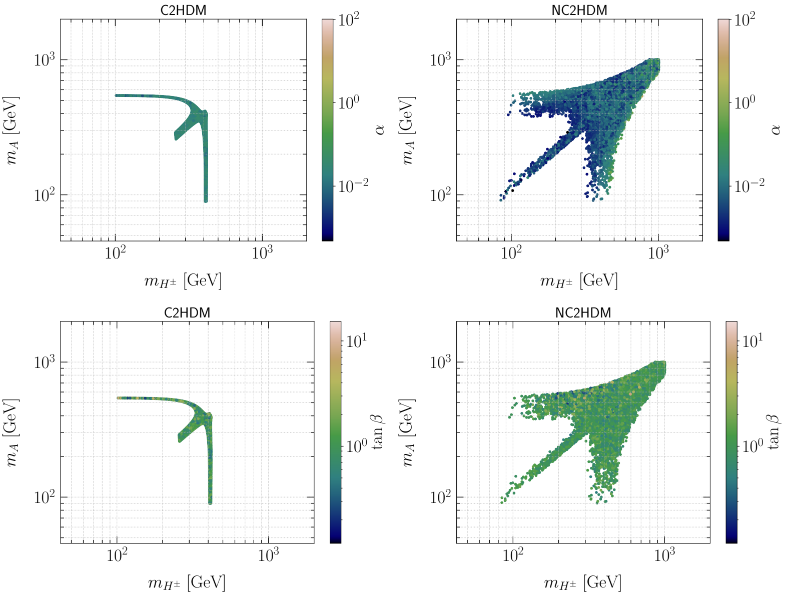
- probability $P(T)$ of being in the false vacuum :

$$P(T) = e^{-I(T)},$$

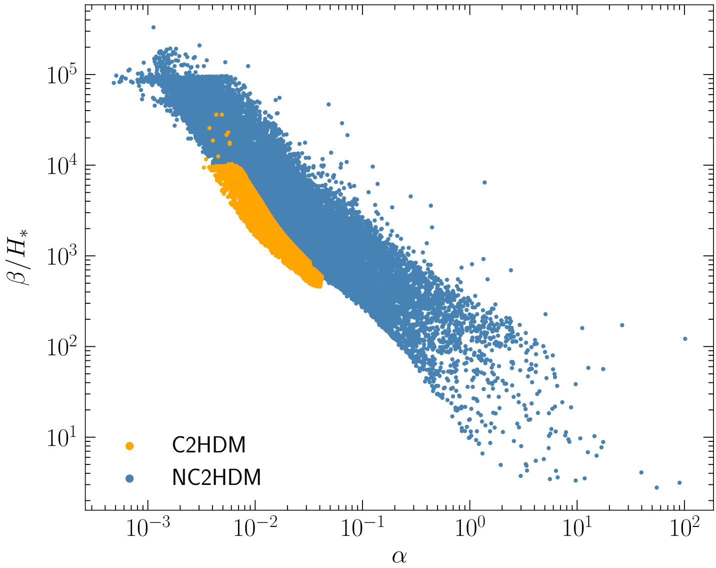
  with $I(T)$, the fraction of the Universe that has transitioned
- percolation temperature T_p:

$$I(T_p) = 0.34$$

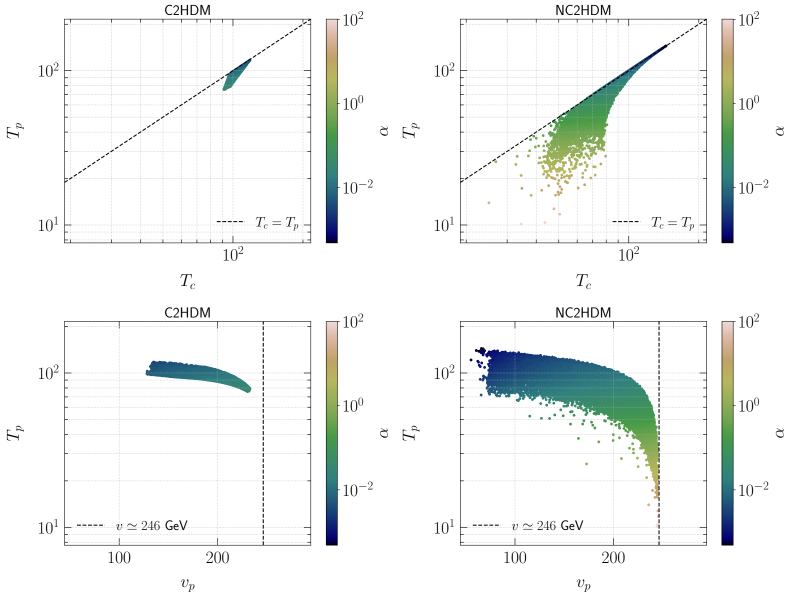
Parameter space



PT parameters: results



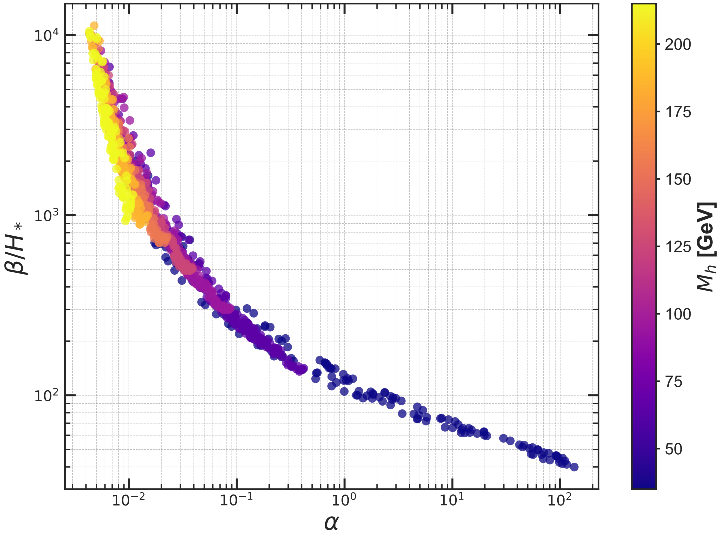
PT parameters: results



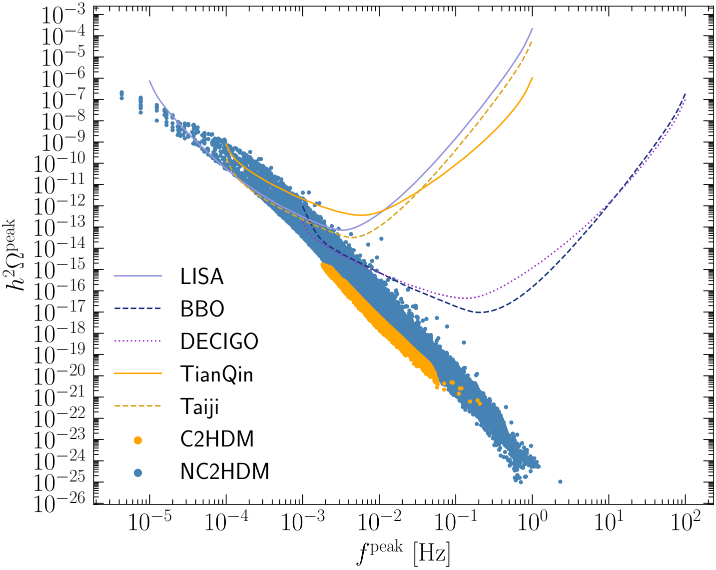
Scalon mass impact on PT strength

M_h [GeV]	T_c [GeV]	T_p [GeV]	α
35	[43.38—59.92]	[5.55—28.46]	[0.246—134.941]
65	[61.39—71.52]	[32.48—54.58]	[0.054—0.401]
95	[78.99—90.22]	[57.11—79.38]	[0.021—0.083]
125	[96.12—108.52]	[78.36—102.10]	[0.011—0.037]
155	[113.00—124.35]	[97.69—120.05]	[0.007—0.021]
185	[129.62—138.40]	[116.16—135.53]	[0.005—0.014]
215	[145.40—150.99]	[133.66—148.31]	[0.004—0.010]

Scalon mass impact on PT strength



Gravitational-wave power spectrum



Outline of the talk

Conclusion

- nearly-conformal dynamics typically lead to significant supercooling because the thermal barrier remains until $T \to 0$
- in the C2HDM however the amount of supercooling is very limited
- NC2HDM offers a larger variability in the amount of supercooling: it can be smaller but also bigger than in the C2HDM case
- a smaller value for the scalon mass (a softer breaking of the scale invariance) leads to stronger FOPT
- C2HDM could not provide a strong enough GW signal to be detected by e.g. LISA, contrary to NC2HDM

The End!

Thank you
for your attention!

Any
questions ?