Higgs emerging from the dark

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We propose a new non-thermal mechanism of dark matter production based on vacuum misalignment. A global X-charge asymmetry is generated at high temperatures, under which both the will-be Higgs and the dark matter are charged. At lower energies, the vacuum changes alignment and breaks the U(1)_X, leading to the emergence of the Higgs and of a fraction of charge asymmetry stored in the stable dark matter relic. This mechanism can be present in a wide variety of models based on vacuum misalignment, and we demonstrate it in a composite Higgs template model, where all the necessary ingredients are naturally present. A light pseudo-scalar is always predicted, with interesting implications for cosmology and future supernova observations.

The presence of Dark Matter (DM) in the Universe is arguably one of the most important mysteries in our knowledge of the physical world. We have compelling evidence in cosmological and astrophysical observations that the majority of the matter density in the whole Universe [1] and around galaxies and galaxy clusters is of non-baryonic nature. Nevertheless, no particle candidate exists within the Standard Model (SM) of particle physics to fill this gap. The most popular paradigm has been the WIMP one, postulating the presence of a new Weakly Interacting Massive Particle beyond the SM. As this paradigm is currently challenged by the non-observation of a signal in direct detection experiments [2], many new mechanisms have been recently proposed: Asymmetric DM [3, 4], freezing-in FIMPs (‘F’ for feebly) [5], 2 → 3 annihilating SIMPs (‘S’ for strongly) [6–8], to name a few.

In this letter we propose a new mechanism for non-thermal DM production based on vacuum misalignment, in models where the Higgs arises as a pseudo-Nambu-Goldstone boson (pNGB) from a spontaneously broken global symmetry. This class of models include composite Higgs models [9], holographic extra dimensions [10, 11], Little Higgs [12, 13], Twin Higgs [14] and elementary Goldstone Higgs models [15]. Our main starting point is the fact that the vacuum of the theory, in general, depends on the temperature of the Universe. Thus, its structure today at zero temperature (where the misaligned Higgs vacuum is needed) and at the global symmetry breaking scale is likely to be different. Within this set-up, we propose that a DM relic density may be asymmetrically produced via a charge that is preserved only in the high-temperature vacuum. At low temperatures, the charge is broken and the fraction of asymmetry stored in Z_2-odd states remain as DM density. The main advantage is that, while the non-thermal DM production is due to an asymmetry, the low energy DM candidate can decouple from the SM thus avoiding conflict with direct detection data. The mechanism we propose requires the following key ingredients:

i) an exact Z_2 symmetry that keeps a DM candidate stable;

ii) at high energies, right below the global symmetry breaking phase transition at T_{HL}, the vacuum of the theory is an essentially Higgsless (HL) phase[1], where the electroweak symmetry is broken at a scale f \gg v_{SM} = 246 \text{ GeV}, and a global U(1)_X symmetry remains unbroken;

iii) at T_{HL}, an asymmetry is generated in - or transferred to - the U(1)_X charged states, some of which are also odd under Z_2;

iv) at T_\ast < T_{HL}, the vacuum starts rotating away from the HL vacuum, and U(1)_X is spontaneously broken;

v) at T \approx 0, the theory settles on the standard pNGB Higgs vacuum, where the misalignment reproduces the electroweak (EW) scale v_{SM}.

We refer to this as a Higgsless phase even though a state with the quantum numbers of the isosinglet Higgs is present because this state is expected very heavy at the transition.
In this process, the fraction of asymmetry stored in $Z_2$-odd states at $T = T_*$ will survive as DM density as long as such states are decoupled from the SM thermal bath. Furthermore, as we will see in an explicit example, the pNGB Higgs emerges from the $Z_2$-even states charged under the $U(1)_X$ while the vacuum rotates away from the HL vacuum.

To demonstrate how the mechanism works, we will consider models of composite Higgs with vacuum misalignment, which can fulfil all the above requirements. For concreteness, we will focus on models based on an underlying gauge-fermion description, for which the symmetry breaking patterns are known \[16, 17\]: the minimal cosets with a Higgs candidate and custodial symmetry are $SU(4)/Sp(4)$ \[18\], $SU(5)/SO(5)$ \[19\] and $SU(4) \times SU(4)/SU(4)$ \[20\]. A $Z_2$ symmetry is already present in the latter case \[21\], while the other two cases can be easily extended to a $SU(6)$ symmetry \[22, 23\]. A global $U(1)_X$ in the HL vacuum has already been used to define a DM candidate in a $SU(4)/Sp(4)$ Technicolor-like theory in Ref. \[24\] (the connection to the composite Goldstone Higgs vacuum has been studied in Ref. \[25\]). We have checked that a $U(1)_X$ can also be defined in the $SU(4) \times SU(4)/SU(4)$ coset (but not in $SU(5)/SO(5)$). Note that the above features can also be found in other cosets that do not have a simple gauge-fermion underlying description, like the models in Refs \[26\, 27\], and can also be found in elementary realisations. Our proposal is therefore rather general.

To better understand the workings of this mechanism, we need to recall some basic information about modern composite Higgs models \[28\, 30\] based on vacuum misalignment: a Higgs-like boson arises as a composite pNGB from the breaking of a global symmetry $G$ to $H$. The model is such that an alignment exists where $H$ contains the EW gauge symmetry $SU(2)_L \times U(1)_Y$. This alignment, however, is not stable as an explicit breaking of $G$ exists in the form of gauge interactions, top couplings to the strong sector, and current masses for the confining fermions. These terms are responsible for generating a vacuum expectation value for the Higgs, which corresponds to a misalignment of the vacuum. We will describe this by an angle, $\sin \theta = v/f$ \[9\]. At $T \approx 0$, we need $v(0) = v_{\text{SM}} = 246$ GeV to reproduce the SM at low energy. The decay constant of the pNGBs (including the Higgs), $f$, is fixed by the degree of fine-tuning in the zero-temperature potential: typically, $\sin \theta \lesssim 0.2$ from electroweak precision measurements \[31\, 33\], thus fixing $f \gtrsim 1.3$ TeV, even though smaller scales may also be allowed \[29\, 35\]. We stress that $f$ is a fixed scale, only depending on the confinement of the underlying strong dynamics, while it’s the value of $v(T)$ at the minimum of the potential that varies with temperature. As we assume that the vacuum is only misaligned along the Higgs direction, the coset structure can be schematically repre-

| pNGBs          | Higgs vacuum $\theta \sim 0$ | HL vacuum $\theta = \pi/2$ | $gX$ |
|----------------|-------------------------------|----------------------------|------|
| $\bar{Q}_0$    | $H_1 = 21/2$                 | $\bar{Q}_X = (h + i\eta)/\sqrt{2}$ | 1   |
| $\bar{H}_0$    | $\eta = 0$                   | $\omega^2, \omega^{3/2}$ | 0   |
| $Z_2$-odd pNGBs| $H_2 = 21/2$                 | $\Theta_1 = -H_0^2 + \frac{2\omega + \omega^{3/2}}{\sqrt{2}}$ | 1   |
| $Z_2$-odd pNGBs| $\Delta = 0$                 | $\Theta_2 = \frac{\omega^{1/2}}{\sqrt{2}}$ | 1   |
| $Z_2$-odd pNGBs| $\varphi = 0$                | $\Theta_1 = \Delta^2 - H^2$ | 1   |
| $Z_2$-odd pNGBs| $\phi = 0$                   | $\Theta_3 = \Delta^2 + H^2$ | 2   |

Note that in Ref. \[22\], the authors focus on a scenario where a $U(1)_{\text{DM}}$ is preserved in the Higgs vacuum, case that is disfavoured by direct detection.

$^2$ Note that in Ref. \[22\], the authors focus on a scenario where a $U(1)_{\text{DM}}$ is preserved in the Higgs vacuum, case that is disfavoured by direct detection.
Assuming $b(T) > 0$, the breaking of the EW symmetry can be achieved for $a(T) > 0$, with the minimum located at $\sin^2 \theta = a(T)/b(T)$ for $0 < a(T)/b(T) < 1$ and at $\sin^2 \theta = 1$ for $a(T)/b(T) \geq 1$. Thus, the vacuum structure needed for our mechanism can be achieved if $a(T)/b(T)$ varies with the temperature and we have:

$$\frac{a(T_{\text{IL}})}{b(T_{\text{IL}})} > 1 \quad \text{and} \quad \frac{a(0)}{b(0)} \ll 1,$$

where $T_{\text{IL}}$ is identified with the temperature of confinement. This implies that the above ratio needs to monotonically decrease with temperature, and that the vacuum is stuck at the IL position until the temperature $T_s$, for which $a(T_s) = b(T_s)$. In this period, for $T_{\text{IL}} > T > T_s$, the electroweak breaking scale is $v(T) = f \gg v_{\text{SM}}$, and the $W$, $Z$ and SM fermions are much heavier than the SM values by a factor $f/v_{\text{SM}}$. If we consider a benchmark scale $f = 1.5$ TeV, this yields $m_{W}^{\text{IL}}(T) = 490$ GeV, $m_{Z}^{\text{IL}}(T) = 560$ GeV, and $m_{\mu}^{\text{IL}}(T) = 1060$ GeV.

We can now start following the thermal history of the DM candidates. At the phase transition temperature, $T_{\text{IL}} \approx O(f)$, the global symmetry $U(1)_X$ is exact while the EW symmetry is broken. The pNGBs, therefore, can be labelled in terms of their electromagnetic and $U(1)_X$ charges. For the template SU(6)/Sp(6) model, refer to the third column of Table I. Note that the will-be Higgs boson forms a $q_X = 1$ state together with the singlet $\eta$, while all the $Z_2$-odd pNGBs have charges $q_X = 1/2$. We will call the latter collectively as $\Theta_i$. One interesting point of our model is that $U(1)_X$, together with baryon and lepton numbers, $B$ and $L$ respectively, has an anomaly with the electroweak gauge interactions, thus if the phase transition at $T_{\text{IL}}$ generates a baryon asymmetry, it will also generate an $X$ asymmetry. The relative densities can be computed following Ref. [23], with the only difference that we will not include any Higgs boson in the computation as our theory is Higgsless at the phase transition. The computation is model-dependent, so here we will show the results for the template model: the relation among the chemical potentials of the various states can be easily computed following their quantum numbers, while we find that the relation imposed by the EW sphalerons is the same as in Ref. [24],

$$2\mu_\Theta + 9\mu_{uL} + 3\mu_W + \mu_L = 0, \quad (4)$$

involving 4 remaining independent chemical potentials: $\mu_\Theta$ of the $Z_2$-odd pNGBs, $\mu_{uL}$ of the left-handed up-type quarks, $\mu_W$ of the $W^-$ boson and $\mu_L = \sum_{i=1}^3 \mu_{L_i}$ being the total one of the three charged leptons. For a strong

1st order phase transition, imposing the vanishing of the total charge and isospin gives a fixed numerical ratio for the asymmetries. Assuming that all $Z_2$-odd pNGBs are light compared to the phase transition scale, we find

$$\frac{X}{B} = -4, \quad \frac{L}{B} = 3/4. \quad (5)$$

We also studied the spectrum of the template model and found that a typical spectrum contains a light pair $\Theta_1 - \Theta_1^*$ (or $\Theta_2 - \Theta_2^*$), while the other two are much heavier; we also find that half of the $X$-charge density is initially stored in $\phi_X$, and the other half in the $Z_2$-odd states $\Theta_i$. For the DM density generated by the asymmetry to persist, it is crucial that the $\Theta_i$ states decouple from the thermal bath before the temperature $T_s$, while the detailed relic density depends on the processes that determine the equilibrium between the $X$-charged states $\phi_X$ and $\Theta_i$. This dynamics, taking place between $T_{\text{IL}}$ and $T_s$ is also very model dependent, however all models have similar qualitative features. We study a simplified scenario where only three states are relevant: $\phi_X$, $\Theta_1$, and $\Theta_1^*$, with $m_{\phi_1} \approx m_{\Theta_1} \approx M_{\phi_X}$ and $f \gg M_{\phi_X} \approx 0$. The latter is justified by the fact that the imaginary part of $\phi_X$, $\eta$, becomes an exact Goldstone at $T = T_s$. The relevant couplings are:

$$\mathcal{L} \supset -\frac{g}{\sqrt{2}} W^+_{\mu} (\Theta_1^* \bar{\Theta}_1^* \sigma^\mu \Theta_1) + \frac{\xi}{2} f \phi_X \Theta_1 \Theta_1 + \text{h.c.} - \frac{g^2}{2} \phi_X^* \phi_X \left( W^+_{\mu} W^-_{\mu} + \frac{1}{2} Z^\mu Z^\mu \right), \quad (6)$$

where $g$ is the SU(2)$_L$ gauge coupling and $\xi$ is a small $U(1)_X$ conserving coupling generated by the pNGB potential. The last term is a relic of the fact that $\phi_X$ contains the will-be Higgs boson, which couples to the massive EW gauge bosons. The coupling $\xi$ is the only one that transfers $X$ charge between the $Z_2$-odd states $\Theta_i$ and $\phi_X$, thus the temperature $T_{dc} = M_{\phi_1}/x_{dc}$, where the two decouple and the $X$ charges in $\Theta_i$ are frozen is determined by this interaction getting off thermal equilibrium: we find that the dominant process is $\Theta_1^* + \phi_X \leftrightarrow \Theta_1 + W^+$, whose cross section can be easily computed. The number density of DM candidates coming from the asymmetry is thus determined by the asymmetry in $\Theta_1$ at $T_{dc}$, which can be computed by solving the appropriate Boltzmann equation. As decoupling

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3 This property has passed largely unnoticed in the literature. Recent application in the context of baryogenesis can be found in Refs. [36, 38].

4 I.e., we assume that the $0^{++}$ state, that may play the role of the Higgs [23], is heavy.

5 For a 2nd order phase transition, the vanishing of the isospin cannot be imposed and there is one remaining free parameter (identifiable with the $L/B$ ratio), so that we predict

$$\frac{X}{B} = -\frac{2(93 + 44\xi)}{63}, \quad \frac{L}{B} = \xi.$$

6 We neglect the couplings of the $Z$, which gives qualitatively very similar results.
occurs where \( \Theta \) are non-relativistic, neglecting the contribution of the W chemical potential, we find:

\[
\frac{\Omega_\Theta}{\Omega_b} \approx \frac{m_\Theta}{m_p} \frac{X}{B} \frac{\Delta n_\Theta(T_{de})}{n_X(T_{de})} \approx \frac{M_\Theta}{1 \text{ GeV}} \times 4 \times e^{-M_\Theta/m_{W}^{\text{HL}}} \times 6 \left( \frac{x_{dc}}{2\eta} \right)^2 e^{-x_{dc}}, \tag{7}
\]

where \( n_X(T_{de}) \) is computed by considering that most of the X asymmetry is stored in \( \phi_X \), which is light (relativistic) and in thermal equilibrium, and \( m_{W}^{\text{HL}} \) is the W mass in the HL vacuum. Assuming that this is the dominant contribution to the DM relic density today, and that \( M_\Theta = M_{DM} \) (the mass of the DM at \( T \approx 0 \) may be different from \( M_\Theta(T_{de}) \)), we can thus solve the decoupling and \( \Omega_\Theta/\Omega_b \approx 5 \) to determine \( x_{dc} \) and \( \xi \). The result is shown in Fig. 1 by the dashed-blue and dashed-red lines, for \( m_{W}^{\text{HL}} = 500 \text{ GeV} \), corresponding to \( f \approx 1.5 \text{ TeV} \). We see that \( \xi \) is required to be very small, and this is a model-building constraint on the explicit models. In composite models the \( \xi \) term breaks the symmetry corresponding to G-parity in QCD and so it is not unreasonable for this to be small. In order for the asymmetry to survive, we need to make sure that \( \Theta \), decouple from the SM thermal bath before the vacuum moves away from the HL vacuum. The freeze-out temperature \( T_F = M_\Theta/x_F \) is determined by the process \( \Theta_1 \Theta_2 \leftrightarrow W^+ W^- \) going out of equilibrium, and the result is shown numerically by the blue solid line in Fig. 1 This implies an upper limit on \( T_* \):

\[
T_* < T_F \approx \frac{M_\Theta}{33} \approx 30 \text{ GeV} \times \frac{M_\Theta}{1 \text{ TeV}}. \tag{8}
\]

As a final consistency check, we determine the freeze-out temperature for \( \phi_X \), coming from the processes \( \phi_X \phi_X^* \leftrightarrow W^+ W^- / ZZ \), giving \( T_F \approx 18 \text{ GeV} \times f/(1.5 \text{ TeV}) \); this temperature is below \( T_F \) for \( M_\Theta \gtrsim m_{W}^{\text{HL}} \). Note that \( \phi_X \) is relativistic at freeze-out.

At the temperature \( T_* < T_F \), the vacuum of the theory starts drifting away from the HL vacuum. At this time, \( U(1)_X \) is spontaneously broken by the vacuum, and \( \phi_X \) is no longer protected from decays. Thus, only the fraction of X charge stored in the \( Z_2 \)-odd pNGBs will survive. The masses of the will-be Higgs \( h \) and of \( \eta \) split, and \( h \) starts acquiring Higgs-like couplings to the SM states, scaling like \( \cos \theta \). Thus, close to the transition temperature \( T_* \), the couplings are still small and the W and Z bosons heavy. While the Universe cools down, the EW masses gradually decrease to the SM values, while the Higgs couplings approach the SM values as \( \cos \theta \to 1 \). In our model, therefore, the 125–GeV Higgs emerges from the dark \( \phi_X \) state during the relaxation at \( T < T_* \). One potential concern is that the \( Z_2 \)-odd states may thermalise again with the SM via the interactions to the W/Z bosons, the top and the emerging Higgs. However, this should not be an issue, as the relaxation away from the HL vacuum can be slow, and the couplings and masses vary slowly with the temperature.

\[\frac{\cos \eta}{f} \left( g^2 \kappa_W W_{\mu
u} \bar{W}^{\mu\nu} + g^2 \kappa_B B_{\mu
u} \bar{B}^{\mu\nu} \right). \tag{9}\]

This term contains a coupling to two photons, proportional to \( \kappa_{\gamma\gamma} = \kappa_W + \kappa_B \), which is very strongly constrained for \( m_\eta \lesssim 1 \text{ GeV} \) (see for instance Ref. [11]), giving rise to bounds on \( f \) many orders of magnitude above the TeV scale. The template model is rather special because it features \( \kappa_{W} = -\kappa_{B} \) so that \( \kappa_{\gamma\gamma} = 0 \) at leading order. \(^7\) Albeit \( \eta \) has a photophobic nature \(^8\), couplings to photons and to SM fermions are generated at loop level \(^{12}\), thus strong bounds may still arise from astrophysics and cosmology. In our case, for \( m_\eta \lesssim 9 \text{ keV} \), strong bounds \( f > O(100) \text{ TeV} \) arise from star evolution \(^{23}\)–\(^{46}\), while for \( m_\eta < 100 \text{ MeV} \), interesting effects...
may be observed in a future supernova observation if \( f \) is in the TeV range \([47]\).

Bounds from cosmology also apply \([48]\), however a detailed analysis is sensitive to the details of the model and of the cosmological evolution of the theory, and they will be presented elsewhere.

In conclusion, in this letter we have presented a new mechanism for non-thermal DM production via vacuum misalignment. The relic density emerges from an asymmetric potential at high energies, while the SM-like Higgs boson also emerges from the high-temperature dark sector. This mechanism predicts a light pNGB from the low-temperature breaking of the U(1) symmetry, leading to interesting effects in future supernova observations and cosmology.

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**[1]** PLANCK collaboration, N. Aghanim et al., “Planck 2018 results. VI. Cosmological parameters,” [\[1807.06209\]].

**[2]** XENON collaboration, E. Aprile et al., “Dark Matter Search Results from a One Ton-Year Exposure of XENON1T,” Phys. Rev. Lett. 121 (2018) 111302, [\[1805.12562\]].

**[3]** S. Nussinov, “Technocosmology: Could a Technibaryon Excess Provide a ‘Natural’ Missing Mass Candidate?,” Phys. Lett. 165B (1985) 55–58.

**[4]** D. E. Kaplan, M. A. Luty and K. M. Zurek, “Asymmetric Dark Matter,” Phys. Rev. D79 (2009) 115016, 0901.4117.

**[5]** L. J. Hall, K. Jedamzik, J. March-Russell and S. M. West, “Freeze-In Production of FIMP Dark Matter,” JHEP 03 (2010) 080, [\[0911.1120\]].

**[6]** E. D. Carlson, M. E. Machacek and L. J. Hall, “Self-interacting dark matter,” Astrophys. J. 398 (1992) 45–52.

**[7]** Y. Hochberg, E. Kuflik, T. Volansky and J. G. Wacker, “Mechanism for Thermal Relic Dark Matter of Strongly Interacting Massive Particles,” Phys. Rev. Lett. 113 (2014) 1402.5143.

**[8]** M. Hansen, K. Langaebke and F. Sannino, “SIMP model at NNLO in chiral perturbation theory,” Phys. Rev. D92 (2015) 075036, 1507.01590.

**[9]** D. B. Kaplan and H. Georgi, “SU(2) x U(1) Breaking by Vacuum Misalignment,” Phys. Lett. 136B (1984) 183–186.

**[10]** R. Contino, Y. Nomura and A. Pomarol, “Higgs as a holographic pseudogoldenbozon,” Nucl. Phys. B671 (2003) 148–174, hep-ph/0306259.

**[11]** Y. Hosotani and M. Mabe, “Higgs boson mass and electroweak-gravity hierarchy from dynamical gauge-Higgs unification in the warped spacetime,” Phys. Lett. B615 (2005) 257–265, hep-ph/0503020.

**[12]** N. Arkani-Hamed, A. G. Cohen and H. Georgi, “Electroweak symmetry breaking from dimensional deconstruction,” Phys. Lett. B513 (2001) 232–240, hep-ph/0105239.

**[13]** N. Arkani-Hamed, A. G. Cohen, E. Katz, A. E. Nelson, T. Gregoire and J. G. Wacker, “The Minimal moose for a little Higgs,” JHEP 08 (2002) 021, hep-ph/0206020.

**[14]** Z. Chatko, H.-S. Goh and R. Harnik, “The Twin Higgs: Natural electroweak breaking from mirror symmetry,” Phys. Rev. Lett. 96 (2006) 231802, hep-ph/0506256.

**[15]** T. Alanne, H. Gertov, F. Sannino and K. Tuominen, “Elementary Goldstone Higgs boson and dark matter,” Phys. Rev. D91 (2015) 095021, 1411.6132.

**[16]** E. Witten, “Current Algebra, Baryons, and Quark Confinement,” Nucl. Phys. B223 (1983) 433–444.

**[17]** D. A. Kosower, “Symmetry breaking patterns in pseudoreal and real gauge theories,” Phys. Lett. 144B (1984) 215–216.

**[18]** J. Galloway, J. A. Evans, M. A. Luty and R. A. Tacchi, “Minimal Conformal Technicolor and Precision Electroweak Tests,” JHEP 10 (2010) 086, 1001.1361.

**[19]** M. J. Dugan, H. Georgi and D. B. Kaplan, “Anatomy of a Composite Higgs Model,” Nucl. Phys. B254 (1985) 299–326.

**[20]** T. Ma and G. Cacciapaglia, “Fundamental Composite 2HDM: SU(N) with 4 flavours,” JHEP 03 (2016) 211, 1508.07014.

**[21]** Y. Wu, T. Ma, B. Zhang and G. Cacciapaglia, “Composite Dark Matter and Higgs,” JHEP 11 (2017) 058, 1703.06903.

**[22]** C. Cai, G. Cacciapaglia and H.-H. Zhang, “Vacuum alignment in a composite 2HDM,” 1805.07619.

**[23]** G. Cacciapaglia, H. Cai, A. Deandrea and A. Kushwaha, “Composite Higgs and Dark Matter Model in SU(6)/SO(6),” 1904.09301.

**[24]** T. A. Ryttov and P. Sannino, “Ultra Minimal Technicolor and its Dark Matter TIMP,” Phys. Rev. D78 (2008) 115010, 0809.0713.
[25] G. Cacciapaglia and F. Sannino, “Fundamental Composite (Goldstone) Higgs Dynamics,” JHEP 04 (2014) 111, 1402.0239.

[26] G. Ballesteros, A. Carmona and M. Chala, “Exceptional Composite Dark Matter,” Eur. Phys. J. C77 (2017) 468, 1704.07388.

[27] R. Balkin, M. Ruhdorfer, E. Salvioni and A. Weiler, “Charged Composite Scalar Dark Matter,” JHEP 11 (2017) 094, 1707.07685.

[28] R. Contino, “The Higgs as a Composite Nambu-Goldstone Boson,” in Physics of the large and the small, TASI 09, proceedings of the Theoretical Advanced Study Institute in Elementary Particle Physics, Boulder, Colorado, USA, 1-26 June 2009, pp. 235–306, 2011. 1005.4269.

[29] B. Bellazzini, C. Csáki and J. Serra, “Composite Higgses,” Eur. Phys. J. C74 (2014) 2766, 1401.2457.

[30] G. Panico and A. Wulzer, “The Composite Nambu-Goldstone Higgs,” Lect. Notes Phys. 913 (2016) pp.1–316, 1506.01961.

[31] K. Agashe, R. Contino, L. Da Rold and A. Pomarol, “A Custodial symmetry for Zb ¯b,” Phys. Lett. B641 (2006) 62–66, hep-ph/0605341.

[32] R. Barbieri, D. Buttazzo, F. Sala, D. M. Straub and A. Tesi, “A 125 GeV composite Higgs boson versus flavour and electroweak precision tests,” JHEP 05 (2013) 069, 1211.5085.

[33] C. Grojean, O. Matsedonskyi and G. Panico, “Light top partners and precision physics,” JHEP 10 (2013) 160, 1306.4655.

[34] D. Ghosh, M. Salvarezza and F. Senia, “Extending the Analysis of Electroweak Precision Constraints in Composite Higgs Models,” Nucl. Phys. B914 (2017) 346–387, 1511.08235.

[35] D. Buarque Franzosi, G. Cacciapaglia and A. Deandrea, “Sigma-assisted natural composite Higgs,” 1809.09146.

[36] S. Bruggisser, B. Von Harling, O. Matsedonskyi and G. Servant, “Baryon Asymmetry from a Composite Higgs Boson,” Phys. Rev. Lett. 121 (2018) 131801, 1803.08546.

[37] S. Bruggisser, B. Von Harling, O. Matsedonskyi and G. Servant, “Electroweak Phase Transition and Baryogenesis in Composite Higgs Models,” JHEP 12 (2018) 099, 1804.07314.

[38] I. Baldo and G. Servant, “High scale electroweak phase transition: baryogenesis & symmetry non-restoration,” JHEP 10 (2018) 053, 1807.08770.

[39] A. Belyaev, M. S. Brown, R. Foardi and M. T. Frandsen, “The Technicolor Higgs in the Light of LHC Data,” Phys. Rev. D90 (2014) 035012, 1309.2097.

[40] D. Liu, J. Low and Z. Yin, “Universal Imprints of a Pseudo-Nambu-Goldstone Higgs Boson,” Phys. Rev. Lett. 121 (2018) 201802, 1805.00489.

[41] M. Bauer, M. Neubert and A. Thamm, “LHC as an Axion Factory: Probing an Axion Explanation for (g − 2), with Exotic Higgs Decays,” Phys. Rev. Lett. 119 (2017) 031802, 1704.08207.

[42] N. Craig, A. Hook and S. Kasko, “The Photophobic ALP,” JHEP 09 (2018) 028, 1805.06538.

[43] M. Bauer, M. Neubert and A. Thamm, “Collider Probes of Axion-Like Particles,” JHEP 12 (2017) 044, 1708.00443.

[44] G. Raffelt and A. Weiss, “Red giant bound on the axion - electron coupling revisited,” Phys. Rev. D51 (1995) 1495–1498, hep-ph/9410205.

[45] A. H. Corsico, O. G. Benvenuto, L. G. Althaus, J. Isern and E. Garcia-Berro, “The Potential of the variable DA white dwarf G117 - B15A as a tool for fundamental physics,” New Astron. 6 (2001) 197–213, astro-ph/0104103.

[46] T. Battich, A. H. Córsico, L. G. Althaus, M. M. Miller Bertolami and M. M. M. Bertolami, “First axion bounds from a pulsating helium-rich white dwarf star,” JCAP 1608 (2016) 002, 1605.07666.

[47] G. G. Raffelt, “Astrophysical axion bounds,” Lect. Notes Phys. 741 (2008) 51–71, hep-ph/0611350.

[48] D. Cadamuro and J. Redondo, “Cosmological bounds on pseudo Nambu-Goldstone bosons,” JCAP 1202 (2012) 032, 1110.2895.