In this article a novel mechanism for dynamical electroweak symmetry breaking and the ensuing appearance of fermion mass terms in the action is proposed. The action contains massless fermions of the SM coupled to gravity through a new type of non-minimal coupling to the vielbein field. The corresponding coupling constants in our approach become zero-dimension scalar fields. Such scalar fields provide the cancellation of the Weyl anomaly \[ \mathcal{A} \].

I. INTRODUCTION

In this article a novel mechanism of dynamical symmetry breaking, thus leading to the appearance of fermion mass terms in the action, is put forward. The assumed action contains the full set of standard model (SM) fermions coupled to gravity through a new type of non-minimal coupling to the vielbein field. In contrast to the fermion-vielbein coupling via a constant term [2], we promote this coupling from a constant to being an actual field in its own right. Correspondingly additional distinct scalar fields are introduced in the theory that furnish a set of coupling coefficients between the various different flavours of SM fermions to the vielbein. Adding scalar fields to the SM action is known to imply the cancellation of the Weyl anomaly [1]. We propose an action made up from the SM fermions coupled to gravity through a new type of non-minimal coupling to the vielbein field. The appearance of fermion mass terms in the action is proposed. The action contains massless fermions of the SM already discussed, while providing sound theoretical predictions for the Higgs mass. In these models the energy scale of the new dynamics was assumed to be at about \( 10^{15} \) GeV. Consistent with this value, these models give naive predictions for the Higgs mass at approximately \( 2m_t \sim 350 \) GeV [26-31], and present experimental data rules out such a value. A smaller value of the Higgs boson mass might be obtained if the calculations rely on large renormalization group corrections [27,30], due to the running of coupling constants between the working scale (\( 10^{15} \) GeV) and the electroweak scale (100 GeV). At any rate this running is not able to explain the observed appearance of the Higgs mass at around 125 GeV. Even more these models require fine-tuning of numerous parameters in order to match experimental results. But that aside, the models with four - fermion interactions should be taken as the phenomenological ones, in which only the leading order in the \( 1/N_c \) expansion is taken into account.

The top see-saw model was advocated initially by Chivukula, Dobrescu, Georgi and Hill [32] in response to the difficulties of the earlier top-quark condensate models. The model admits, in addition to the top quark, an additional heavy fermion, \( \chi \). In particular in the top see-saw model there is no need for fine tuning of parameters, and a later version of the model by Dobrescu and Cheng [34] includes a light composite Higgs boson.

In the 1970s the technicolor (TC) model was formulated as an attempt to resolve the shortcomings of the SM already discussed, while providing sound theoretical predictions for the Higgs mass. TC is a gauge
theory of fermions with no elementary scalars, yet at
the same time inclusive of dynamical electroweak sym-
metry breaking and flavor symmetry breaking. In the
TC model the Higgs, instead of being a fundamental
scalar is composed of a new class of fermions called
technifermions interacting via technicolor gauge bosons.
This interaction is attractive and as a result, by anal-
gogy with Bardeen-Cooper-Schrieffer (BCS) superconduc-
tor theory, leads to the formation of fermion conden-
sates. However, TC theory fails to produce mass terms
for the quarks and leptons. For this reason the model
was extended to the Extended Technicolor (ETC) model
[37, 38], but this model is inconsistent with experimental
constraints on flavor changing neutral currents, and pre-
cision electroweak measurements. Later on an improved
TC model called the walking technicolor model [39, 40]
was constructed that resolves these issues, although it
still fails to predict correctly the measured value of the
top-quark mass. For a detailed pedagogical review of the
TC and ETC models see Refs. [24–41].

After the technicolor models came the idea that the
Higgs boson appears as a pseudo-Goldstone boson, as
noted originally in Ref. [72]. This idea forms the basis
of Little Higgs Models [43–45]. The model stems a well
known result of Goldstone’s theorem, namely that sponta-
naneous breaking of a global symmetry yields massless
scalar particles, or Goldstone bosons. By selecting the
right global symmetry, it is possible to have Goldstone
bosons that correspond to the Higgs doublet in the SM.

Subsequently the partial compositeness model was con-
j ectured, by Kaplan [49] initially. Therein each SM par-
eticle has a heavy partner that can mix with it. Like this
the SM particles are linear combinations of elementary
and composite states via a mixing angle. The elegance of
the partial compositeness model is its simplicity, without
the need for deviations beyond the standard model, thus
relying on the known fundamental particles only [50].

More recently, among a variety of theoretical papers
that appeared between December 2015 and August 2016,
there are several that consider both the 125 GeV Higgs
boson, and the hypothetical new heavier Higgs bosons as
composite due to the new strong interaction, as noticed
for example in Refs. [51–52]. In a different approach
there are other papers devoted to the description of the
composite nature of the heavier (750 GeV) Higgs boson
only [53–60].

Other more modern composite Higgs models consist of
the walking model [67, 68], the ideal walking model [69–
72], the technicolor scalar model [73], Sannino’s model
of the generalized orientifold gauge theory approach to
electroweak symmetry breaking [74], and the walking model
in higher-dimensional $SU(N)$ gauge theories of Dietrich
and Sannino [75, 76]. A more comprehensive review of
modern composite Higgs models can be found in [77].

Quantum gravity in the first order formalism with ei-
ther the Palatini action [78, 79] or the Holst action [80],
shares a property with the above composite Higgs mod-
els: it leads to a four - fermion interaction. This four
fermion interaction is now between spinor fields coupled
in a minimal way to the torsion field [81, 82], but im-
portantly this four - fermion interaction leads to fermion
condensates of the type that could constitute the com-
posite Higgs. A similar idea has already been suggested
in Ref. [83], namely that the torsion field coupled in a
non-minimal way to fermion fields [84–86] could be a
source of dynamical electroweak symmetry breaking.

This type of action with fermion fields coupled to the
vielbein is one of the main ingredients of our action. The
other main ingredient flows from a theory that contains
fundamental scalar fields with zero mass-dimension re-
cently proposed in [1], as a way of simultaneously can-
celing both the vacuum energy and the Weyl anomaly.
Starting with a theory of the SM coupled to gravity, the
fermion and gauge fields are coupled to a background
classical gravitational field, with a right-handed neutrino
for each generation also coupled to the background field.
This representation is a QFT on a classical background
spacetime. The Weyl anomaly is a measure of the failure of
the classically Weyl-invariant theory to define a Weyl-
invariant quantum theory, and can be expressed in terms
of the number of different fields present in the theory.
In the SM there are $n_0 = 4$ ordinary real scalars in the usual
complex Higgs doublet, $n_{1/2} = 3 \times 16 = 48$ Weyl spinors
(16 per generation), $n_1 = 8 + 3 + 1 = 12$ gauge fields
of $SU(3) \times SU(2) \times U(1)$, and a single gravitational field
($n_2 = 1$). With these combinations the Weyl anomaly is
non-zero. However suppose that the Higgs and graviton
fields are not fundamental fields but rather composite,
such that their contributions to the vacuum energy and
Weyl anomaly can be dropped, and $n_0 = n_2 = 0$. The im-
plication is that by introducing $n_0' = 36$ scalars with zero
mass-dimension, then not only does the Weyl anomaly
vanish, but also the vacuum energy vanishes as well.

This embodies our motivation for the starting point of
this work: a model comprising Dirac fermions of the SM
with a non-minimal coupling to gravity, that includes 36
fundamental scalar fields. The details of the construc-
tion of the theory is described in full in §11A. We find
that in this model, the theory admits mass terms for the
fermions. By using the Schwinger-Dyson equation for the
fermion self-energy function, a relationship is derived
between the top-quark mass, the Planck mass (the ultra-
violet cut-off point of the loop integral) and the strength
of the inverse coupling between the fermion fields and the
vielbein. From this relation an estimation for the mass
of the top-quark can be extracted.

This paper is organized as follows. In Section II the
first main ingredient of the action is constructed, the
piece that includes Dirac fermions coupled non-minimally
to the vielbein field that describes the background gravi-
tational field. In §11 the remaining piece of the action is
formulated, namely the dynamical piece of the fundamen-
tal scalar fields described above. In §IV the Schwinger-
Dyson approach is implemented to calculate the leading-
order piece of the self-energy of the dominant scalar field
in the action, namely the scalar field that couples to the
top-quark, being the heaviest flavour. In this section the main result of the paper is derived, namely the emergence of a mass term for the top quark. A relationship between the top-quark mass, the coupling of the scalars to the vielbein (gravitational) field, and the cut-off energy of the loop integral (the Planck mass) is used to estimate the leading-order contribution to the top-quark mass. In §VII the main conclusions of the paper are summarized along with prospects for future research. There are two appendices: appendix §A contains the main formulas for a Wick rotation used in the computation of the one-loop integral, and appendix §B contains full details of the angular integrals in the one-loop integral.

Throughout this article lower case latin letters a, b, c, d = 0, 1, 2, 3 label internal Lorentz indices, greek letters μ, ν, … = 0, 1, 2, 3 label spacetime indices, and εabcd is the completely antisymmetric Levi-Civita symbol. Lorentz indices are raised and lowered by the Minkowski metric ηab and spacetime indices are raised and lowered by a spacetime metric gμν. Metrics are assumed to have a mostly negative signature, and specifically the Minkowski metric is ηab = diag(1, −1, −1, −1).

II. FUNDAMENTAL SCALARS IN MODELS WITH THREE GENERATIONS OF SM FERMIONS

A. Non-minimal coupling of fermions with gravity

The action for Dirac fermions coupled to the vielbein field can be written as [81, 87]

\[ S = \frac{1}{6} \epsilon_{abcd} \int d^4x \theta^a \wedge e^b \wedge e^c \wedge e^d, \]

where

\[ \theta^a = \frac{i}{2} \left[ \bar{\psi} \gamma^a D_\mu \psi - D_\mu \bar{\psi} \gamma^a \psi \right] dx^\mu, \]

with γa the usual Dirac matrices, \( \bar{\psi} = \psi^\dagger \gamma^0 \), and \( D_\mu \psi \) is the covariant derivative acting on the Dirac fermion field defined by [88]

\[ D_\mu \psi = \partial_\mu \psi - \frac{i}{2} \gamma_{ab} \omega^{ab}_\mu \psi, \]

where γab = ½ [γa, γb] and \( \omega^{ab}_\mu \) is the spin connection. The fields e^a are the vielbein fields that describe the background gravitational field, related to the spacetime metric g through the completion relation \( \eta_{ab} e^a_\mu e^b_\nu = g_{\mu\nu} \). The action in (1) is a functional of e alone, as opposed to being a functional of both e and ω. Explicitly, ω becomes dependent on e by imposing Cartan’s first structure equation.

In turn, the action for the right-handed Weyl fermions may be written as Eq. (1) with \( \theta^a \rightarrow \theta_R^a \), where

\[ \theta_R^a = \frac{i}{2} \left[ \bar{\psi} R \sigma^a D_\mu \psi_R - D_\mu \bar{\psi}_R \sigma^a \psi_R \right] dx^\mu, \]

while the action for the left-handed Weyl fermions may be written as Eq. (1) with \( \theta^a \rightarrow \theta_L^a \) and

\[ \theta_L^a = \frac{i}{2} \left[ \bar{\psi}_L \sigma^a D_\mu \psi_L - D_\mu \bar{\psi}_L \sigma^a \psi_L \right] dx^\mu, \]

A non-minimal coupling of fermions to gravity is achieved via the substitution [2]

\[ \theta_R^a \rightarrow \theta^a + \frac{i}{2} \left[ \bar{\psi}_R \xi_R \sigma^a D_\mu \psi_R - \xi_R \overline{D_\mu \psi_R} \right] dx^\mu \]

and

\[ \theta_L^a \rightarrow \theta^a + \frac{i}{2} \left[ \bar{\psi}_L \xi_L \sigma^a D_\mu \psi_L - \xi_L \overline{D_\mu \psi_L} \right] dx^\mu \]

with complex-valued constants \( \xi_R \) and \( \xi_L \). Notice that in the absence of the spin connection, \( D_\mu = \partial_\mu \) and the imaginary parts of \( \xi_R \) and \( \xi_L \) decouple and do not interact with fermions. At the same time the real parts of these constants may be absorbed by a rescaling of the fermion fields.

That said, suppose that this construction is extended to the case where \( \xi_R \) and \( \xi_L \) become coordinate-dependent fields. In this case both the real and imaginary parts of \( \xi_R \) and \( \xi_L \) interact with the fermions of the theory, and cannot be removed by any rescaling of the fields. Furthermore, let \( \xi_R \) and \( \xi_L \) be assigned flavour indices. In the presence of SU(3) ⊗ SU(2) ⊗ U(1) gauge fields, there are \( n_1 = 8 + 3 + 1 = 12 \) vector fields. In accordance the number of Weyl fermions in the SM with three generations: \( n_1/2 = 3 \) leptons + 3 quarks \( \times 2_{\text{up} \& \text{down}} \times 2_{\text{left} \& \text{right}} \times 3_{\text{generations}} = 48 \).

B. Parity breaking interactions with zero dimension scalar fields

The theory can be extended one stage further to the case where \( \xi_L \) and \( \xi_R \) carry not only flavour indices but also generation indices. Even more they can be further distinguished by assigning different forms for the matrices \( \xi \) for the left-handed and the right-handed particles, while \( \xi \) the matrices for the quarks and leptons remain identical. In this approach the fermion term \( \theta^a \) in (1) now reads

\[ \theta^a = \theta^a_L + \theta^a_R, \]

where

\[ \theta_R^a = \frac{i}{2} \left[ \bar{\psi}_R (1 + \xi_R) \sigma^a D_\mu \psi_R \right. \]

\[ \left. - (D_\mu \bar{\psi}_R) (1 + \xi_R^\dagger) \sigma^a \psi_R \right] dx^\mu, \]

\[ \theta_L^a = \frac{i}{2} \left[ \bar{\psi}_L \sigma^a D_\mu \psi_L - \xi_L \overline{D_\mu \psi_L} \right] dx^\mu, \]

\[ \theta_L = \theta_L \rightarrow \theta_L^\dagger \]

\[ \theta_R = \theta_R \rightarrow \theta_R^\dagger \]
There would be the same number of scalar fields as above needed for the cancellation of the Weyl anomaly [1]. This gives rise to \( n_0' = 2 \times 2 \times 3 \times 3 = 36 \) components of the scalar fields. This is precisely the number of scalars \( n_0' = 3n_1 \) needed for the cancellation of the Weyl anomaly [1].

C. Interactions that conserve parity

There would be the same number of scalar fields as above if \( \xi_L \) were identical to \( \xi_R \), but with the matrices \( \xi \) being different for quarks and leptons. In this case the interaction with \( \xi \) does not break parity. The fermion action is still defined as in Eq. (1), but instead of Eq. (8),

\[
\theta^a = \theta^a_1 + \theta^a_q,
\]

where

\[
\theta^a_1 = \frac{i}{2} \left[ \bar{\psi}_q (1 + \xi) \gamma^a D_\mu \psi_q - (D_\mu \bar{\psi}_q) (1 + \xi^\dagger) \gamma^a \psi_q \right] dx^\mu ,
\]

and

\[
\theta^a_q = \frac{i}{2} \left[ \bar{\psi}_q (1 + \xi) \gamma^a D_\mu \psi_l - (D_\mu \bar{\psi}_q) (1 + \xi^\dagger) \gamma^a \psi_l \right] dx^\mu .
\]

Here both \( \xi_q \) and \( \xi_l \) are the \( 3 \times 3 \) complex-valued matrices in flavor space. As before, \( n_0' = 2 \times 3 \times 3 = 36 \) zero dimension scalar fields.

D. \( \Phi \) - fields

It is worth mentioning that \( \theta \) - field of Eq. (12) may be replaced in a different way in order to produce the zero dimension scalar fields:

\[
\theta^a \rightarrow \frac{i}{2} \left[ \bar{\psi} \gamma^a D_\mu \psi - (D_\mu \bar{\psi}) \gamma^a \psi \right] dx^\mu + \frac{i}{2} \left[ \bar{\psi}^A \Phi^a_{AB} D_\mu \psi^B \right] dx^\mu .
\]

Here \( \Phi \) is the \( SU(2) \) doublet that carries in addition the generation index \( A \) and the flavor \( B \) of the right-handed fermions. (We may take \( e, \mu, \tau \) for leptons, and \( d, s, b \) and/or \( u, c, t \) for quarks.) This expression is the further extension of the non-minimal coupling of fermions to gravity, in which the above-mentioned constants \( \xi \) not only become fields, but are, in addition, transformed under the \( SU(2)_L \) group. For example, we can take \( A = 1, 2, 3 \) corresponding to the three generations of quarks, and \( B = 1, 2, 3 \) corresponding to the \( u, c, t \) quarks. As a result we obtain \( 4 \times 4 \times 3 = 144 \) real-valued components of these scalar fields. This number is larger than the one needed for the cancellation of the Weyl anomaly [1]. The minimal choice here is if we consider the singlets in flavor space, which gives only \( 4 = 16 \) components of the scalars - the number that is smaller than the required 32 components. This scheme, therefore, cannot be used directly to build the theory with the cancellation of Weyl anomaly. The other ingredients should be added in order to match this condition. Therefore, in the following we will concentrate on the scheme proposed above in Sect. II C.

III. ELECTROWEAK SYMMETRY BREAKING AND MASS GENERATION

Consider the SM with three generations and the fundamental scalar \( \xi \) fields constructed in II C. With variations of the vielbein and spin connection restrained, define an action for the zero dimension scalar fields as

\[
S_B = \alpha_{ABCD}^{uv} \int d^4x \left(\xi_u^{AB}\right)^f \square^2 \xi^{CD}_\epsilon ,
\]

where \( A, B, C, D = 1, 2, 3 \) are generation indices, \( u, v = q, l \) are flavor indices, and \( \alpha_{ABCD}^{uv} \) are a set of coupling constants. The effective four - fermion interactions arise from the exchange by quanta of the fields \( \xi_q \) and \( \xi_l \). The effective four - fermion interaction is non-local.

Consider now a certain sector of the theory that describes the dominant contributions of the field \( \xi^{33}_q \). Add this piece of the action to the action in (11), with \( \theta^a \) given by (11), to obtain

\[
S = \int d^4x \left[ \alpha \xi \square^2 \xi + \frac{i}{2} Q (1 + \xi) \bar{D} Q - \frac{i}{2} \bar{D} Q (1 + \xi^\dagger) Q \right]
\]

where \( \bar{D} = \gamma^\mu D_\mu, \xi = \xi^{33}_q, \alpha = \alpha^q_{3333} \) and

\[
Q_3(x) = \begin{pmatrix} t \\ b \end{pmatrix}
\]

is the third generation of quarks. Note that the covariant derivative \( D \) contains gauge fields. Define an effective action, \( S_4 \) as \( e^{i S_4} = \frac{1}{4 \alpha} \int D x \xi \frac{1}{2} \bar{D} \xi e^{i S} \). After integrating out the scalar fields the corresponding effective action is found to be

\[
S_4 = -\frac{1}{4 \alpha} \int d^4x d^4y \bar{Q}_3(x) \gamma^\mu D_\mu Q_3(x) \square^{-2}(x, y) D_\mu \bar{Q}_3(y) \gamma^\nu Q_3(y),
\]
where \( \Box^{-2}(x, y) \) is the square of the inverse propagator of the boson field \( \xi^3 \). The one-loop contribution to the two-point Green function is straightforwardly read off Eq. (15). The form of the self-energy of the third generation quarks then follows, as written in Eq. (19) below. The corresponding Feynman diagram is shown in Fig. 1.

![Feynman Diagram](image)

**FIG. 1:** Feynman diagram of the fermion self-energy corresponding to Eq. (19). The incoming and outgoing lines are initial and final fermion states with momentum \( p \). The dashed line in the loop is a scalar boson propagator carrying virtual momentum \( p - k \), and the thick horizontal line corresponds to the full fermion propagator with renormalized mass \( \Sigma(k) \) inclusive of all self-energy corrections, as related in Eq. (19).

### IV. SCHWINGER-DYSON APPROACH FOR CALCULATING FERMION MASS

At this point the discussion follows the method of [89], itself based on [90]. The idea of this method is to write down the Schwinger-Dyson equation in order express the inverse of the full propagator, \( D^{-1}(p) \) in terms of the self-energy function \( \Sigma(p) \) in the form of a simplified ansatz, through which \( \Sigma(p) \) can be determined. The self-energy function, \( \Sigma(p) \) of the third generation quarks through leading order, has the form

\[
\Sigma(p) = \frac{1}{\alpha} \int d^4k \frac{\gamma k - \Sigma(k)}{\gamma k - \Sigma(k)} \frac{1}{(p - k)^4}, \tag{19}
\]

where the standard notation \( \int d^4k (\ldots) = (2\pi)^{-4} \int d^4k (\ldots) \) has been used. Now apply a Wick rotation (see Appendix A for details) to obtain

\[
i\Sigma(p) = \frac{1}{\alpha} \int d^4k E \frac{1}{\gamma E k E - i\Sigma(k)} \frac{1}{(p_E - k_E)^4}. \tag{20}
\]

Since all expressions from now on are in terms of Euclidean-space variables the subscript ‘\( E \)’ can be dropped.

The ansatz for the Schwinger-Dyson equation has the form

\[
D^{-1}(p) = A(p^2)\gamma p - iB(p^2) = \gamma p - i\Sigma(p). \tag{21}
\]

The right-hand side of (21) may be substituted inside (20). To keep notation brief it is useful to use the shorthand \( A = A(k^2) \) and \( B = B(k^2) \). The resulting expression is

\[
i\Sigma(p) = \frac{1}{\alpha} \int d^4k \frac{\gamma k}{(A^2k^2 + B^2)(p - k)^2} \left[1 - \frac{1}{(p - k)^4}\right]. \tag{22}
\]

By (A3), \( (\gamma k)^2 = \gamma^\mu \gamma^\nu k_\mu k_\nu = \frac{1}{2} \delta^{\mu\nu} k_\mu k_\nu = k^2 \), and hence

\[
i\Sigma(p) = \Sigma_1(p) + \Sigma_2(p), \tag{23}
\]

where

\[
\Sigma_1(p) := \frac{1}{\alpha} \int d^4k \frac{k^2 A}{(A^2k^2 + B^2)(p - k)^2} \frac{1}{(p - k)^4}. \tag{24}
\]

\[
\Sigma_2(p) := \frac{i}{\alpha} \int d^4k \frac{k^2 B}{(A^2k^2 + B^2)(p - k)^2} \frac{1}{(p - k)^4}. \tag{25}
\]

\( \Sigma_1(p) \) needs to be in a form in which \( \gamma p \) appears explicitly such that \( A(p^2) \) can be easily read off (21) by comparing coefficients. Multiply \( \Sigma_1 \) by \( \gamma p \) to obtain

\[
\frac{1}{4} \text{Tr} \gamma p \Sigma_1(p) = -\frac{1}{2\alpha} \int d^4k \frac{k^2 A}{(A^2k^2 + B^2)} \left[1 - \frac{1}{(p - k)^2} - \frac{p^2 + k^2}{(p - k)^4}\right]. \tag{26}
\]

Now multiply (26) by \( \gamma p \), bearing in mind that by (A3) \( (\gamma p)^2 = p^2 \), to obtain

\[
p^2 \Sigma_1(p) = -\frac{1}{2\alpha} \gamma p \int d^4k \frac{k^2 A}{(A^2k^2 + B^2)} \left[1 - \frac{1}{(p - k)^2} - \frac{p^2 + k^2}{(p - k)^4}\right]. \tag{27}
\]

In this form the \( \gamma p \) term is explicit and (23) reads

\[
i\Sigma(p) = \frac{1}{2\alpha p^2} \gamma p \int d^4k \frac{k^2 A}{(A^2k^2 + B^2)} \left[1 - \frac{1}{(p - k)^2} - \frac{p^2 + k^2}{(p - k)^4}\right] + \frac{i}{\alpha} \int d^4k \frac{k^2 B}{(A^2k^2 + B^2)} \frac{1}{(p - k)^2}. \tag{28}
\]

Substitute (28) back into (21) and compare coefficients on both sides to obtain the integral equations

\[
A(p^2) = 1 + \frac{1}{2\alpha p^2} \int d^4k \frac{k^2 A}{(A^2k^2 + B^2)} \left[1 - \frac{1}{(p - k)^2} - \frac{p^2 + k^2}{(p - k)^4}\right], \tag{29}
\]

\[
B(p^2) = \frac{1}{\alpha} \int d^4k \frac{k^2 B}{(A^2k^2 + B^2)} \frac{1}{(p - k)^2}. \tag{30}
\]

Write (29) and (30) in terms of four-dimensional polar coordinates:
\[
A(p^2) = 1 + \frac{1}{2\alpha p^2} \frac{1}{(2\pi)^2} \int_0^\infty dk \int_0^\pi d\theta_1 \int_0^\pi d\theta_2 \int_0^{2\pi} d\theta_3 k^3 \sin^2 \theta_1 \sin \theta_2 \frac{k^2 A}{(A^2 k^2 + B^2)} \left( \frac{1}{p^2 + k^2 - 2pk \cos \theta_1} - \frac{p^2 + k^2}{(p^2 + k^2 - 2pk \cos \theta_1)^2} \right)
= 1 + \frac{\pi}{\alpha p^2} \frac{1}{(2\pi)^2} \int_0^\infty dk \int_0^\pi d\theta_1 \frac{k^4 A}{(A^2 k^2 + B^2)} \left( \frac{1}{p^2 + k^2 - 2pk \cos \theta_1} - \frac{p^2 + k^2}{(p^2 + k^2 - 2pk \cos \theta_1)^2} \right) \quad (31)
\]

\[
B(p^2) = \frac{1}{\alpha} \frac{1}{(2\pi)^4} \int_0^\infty dk \int_0^\pi d\theta_1 \int_0^\pi d\theta_2 \int_0^{2\pi} d\theta_3 \sin^2 \theta_1 \sin \theta_2 \frac{k^2 B}{(A^2 k^2 + B^2)} \left( \frac{1}{p^2 + k^2 - 2pk \cos \theta_1} - \frac{p^2 + k^2}{(p^2 + k^2 - 2pk \cos \theta_1)^2} \right) \quad (32)
\]

The integrals in (37) and (38) are ultraviolet divergent. However, in the integrals, (39) and (39), except for the first term (independent of \( p \)) are infra-red divergent. However, in
our scheme we implement an infra-red cutoff in addition to an ultraviolet cutoff. Physically the presence of this infra-red cutoff is needed in order to satisfy present experimental bounds on the existence of extra scalar excitations. From this bound we derive that the infra-red cutoff in the above integral should be of the order \( \lambda \sim 1 \) TeV. Evidently, the dominant contribution to \( B(p^2) \) originates from the large \( k \) region, provided that \( \lambda \gg m_t \). The implication is that \( B \) is log divergent. The integral over \( k^2 \) is assumed to have an upper limit at \( k = \Lambda \) with \( \Lambda \) finite. This essentially means that the four momentum in the loop is bound from above (see Fig. 1). The upshot is that \( \Lambda \) is large to the extent that the value of the integral is saturated by the region where \( k \) is close to \( \Lambda \).

In physical terms the loop is dominated by large values of \( k \), or equivalently, the loop is dominated by small distances of the order \( 1/\Lambda \), with \( \Lambda \) of the order of the Planck mass. As a result it is reasonable to take \( p \) to be of the order of 100 GeV, close to the fermion (top-quark) mass. Since 1 GeV is much smaller than the Planck scale, \( p \) is effectively zero in the integral.

The leading order in \( p^2 \) contributions to \( A \) and \( B \) are straightforwardly read off Eqs. (39) and (40). By inserting a cut-off at the ultra-violet and at the infra-red end of the spectrum, these LO contributions have the forms

\[
A_0 = 1 - \frac{1}{16\pi^2\alpha} \int_{\lambda^2}^{A^2} \frac{dk^2}{A_0^2(k^2 + B_0^2)} ,
\]

\[
B_0 = \frac{1}{16\pi^2\alpha} \int_{\lambda^2}^{A^2} \frac{dk^2}{A_0^2(k^2 + B_0^2)} .
\]

By canceling a constant factor of \( B_0 \) on both sides of (42) and substituting a reparameterization of the form \( B_0 = m_t A_0 \), Eqs. (41) and (42) reduce to

\[
A_0 = 1 - \frac{1}{16\pi^2\alpha} \int_{\lambda^2}^{A^2} \frac{dk^2}{A_0^2(k^2 + m_t^2)} ,
\]

and

\[
1 = \frac{1}{16\pi^2\alpha} \int_{\lambda^2}^{A^2} \frac{dk^2}{A_0^2(\lambda^2 + m_t^2)} ,
\]

respectively. Substitute (44) in (43) to obtain

\[
A_0 = 1 - A_0 \quad \Rightarrow \quad A_0 = \frac{1}{2}.
\]

Now substitute (45) in (44) to find

\[
1 = \frac{1}{4\pi^2\alpha} \ln \left( \frac{A^2 + m_t^2}{\lambda^2 + m_t^2} \right) .
\]

This gives rise to the existence of a critical value of the coupling constant

\[
\alpha_c = \frac{1}{4\pi^2} \ln \left( \frac{A^2}{\lambda^2} \right) \sim 1.86641 .
\]
V. CONCLUSION

In this paper we have proposed a novel scheme that could pave the way for the construction of the ultra-violet completion of the SM. The said mechanism, instead of the conventional Higgs boson field, features a set of 36 zero-dimension scalar fields that give rise to dynamical symmetry breaking. In consequence we have presented a mechanism for obtaining fermion mass terms in the SM action. We have implemented this mechanism to the top-quark sector of the action, and shown that it is capable of predicting a value for the top quark mass in agreement with the measured value, within specific limits imposed on the energy scale of the interactions between the scalars and the top-quark.

These scalar fields are precisely the couplings between the fermions and the vielbein field of the background gravitational field. The form of the coupling we propose is a direct extension of the non-minimal coupling of fermions to the vielbein \[ \gamma \], in which the constant coupling terms are promoted to fields in their own right. We further generalize to a scenario where the coupling terms between the vielbein and the various different fermion flavours are distinctly different fields for each different fermion in the SM action. The outcome is that we obtain a total of 36 scalar fields that constitute the coupling terms. At the same time this is the precise number of zero-dimension scalar fields required to have a vanishing Weyl anomaly \[ \gamma \], this being the case for the action we have formulated.

But more fundamentally, in the model that we propose the zero-dimension scalar fields have manifestly geometrical origins. We assume that the sector of the theory containing these zero-dimension scalar fields, is valid within the range of energy scales between \( \lambda \approx 1 \, \text{TeV} \) and \( \Lambda \approx m_p \sim 10^{19} \, \text{GeV} \). The presence of the infrared cutoff, \( \lambda \) is necessary in order to satisfy the currently known lower bound on extra scalar fields in the SM. The ultraviolet cutoff, \( \Lambda \) is naturally of the order of the Planck mass, the scale at which quantum gravity effects are expected to play a significant role.

On grounds that the Weyl anomaly vanishes in the theory presented in this work, then arguably this theory offers a solution for the cosmological constant problem. We argue that within this theory the fermion masses are generated dynamically due to the exchange of such scalars between the SM fermions. We have tested our model by implementing a simplified toy-model version, containing just a doublet of the top quark and bottom quark interacting via the exchange of the zero-dimension scalars. Specifically, we used an approach based on the truncated Schwinger - Dyson equations (the so - called rainbow approximation), which enables a description of dynamical symmetry breaking. Our observations show that there exists a certain critical value of the coupling constant \( \alpha \). For smaller values (strong coupling) the dynamical fermion mass is generated. In order for the fermion mass to be close to the top quark mass, the value of \( \alpha \) must be close to the critical value.

Thinking beyond the results in this article, the toy model discussed here yields equal masses for both the top and bottom quarks. Without doubt this model should be generalized to incorporate different masses for the top and bottom quarks. In particular a prospective generalized model should have a suppression of the bottom-quark mass in comparison with the top quark mass. Above all, the interactions between other components of the scalars should be taken into account, which give rise to the smaller masses of these fermions. Further afield the question arises of how to incorporate the generation of the difference in masses between the \( u \) and \( d \) quarks, the \( c \) and \( s \) quarks, as well as the difference in masses between the charged leptons and their neutrinos. This generalized theory is beyond the scope of this paper but will be addressed in a follow up paper.

ACKNOWLEDGMENTS

This work has been supported by the European Research Council (ERC) under the European Union’s Horizon 2020 research and innovation programme (Grant Agreement No. 694248).

Appendix A: Wick rotations

In Minkowski space the \( \gamma \) matrices satisfy the Clifford algebra given by

\[
\{ \gamma^\mu, \gamma^\nu \} = 2\eta^{\mu\nu}, \quad (\mu, \nu = 0, 1, 2, 3), \tag{A1}
\]

where \( \eta^{\mu\nu} = \text{diag}(1, -1, -1, -1) \). A Wick rotation comprises a transformation of the \( \gamma \) matrices to new matrices, \( \gamma_E \) defined by

\[
\gamma_E^A = \gamma^0, \quad \gamma_E^i = -i\gamma^i, \quad (i = 1, 2, 3) \tag{A2}
\]

such that

\[
\{ \gamma_E^\mu, \gamma_E^\nu \} = 2\delta^{\mu\nu}, \quad (\mu, \nu = 1, 2, 3, 4). \tag{A3}
\]
Four-vector components get transformed under a Wick rotation to new components. For example given components $k^\mu$ of the four vector $k$ in Minkowski space,

$$k^x_E = ik^0, \quad k^z_E = k^1, \quad (i = 1, 2, 3). \quad \text{(A4)}$$

The implication is that

$$k^2 = (k^0)^2 - (k^1)^2 = -(k^1_E)^2 - (k^2_E)^2 - (k^3_E)^2 \equiv -k^2_E. \quad \text{(A5)}$$

$$\gamma k = \gamma^0 k^0 - \gamma^i k_i = -i\gamma^i_E k^i_E - i\gamma^1_E k^1_E - i\gamma^2_E k^2_E - i\gamma^3_E k^3_E \equiv -i\gamma_E k_E. \quad \text{(A6)}$$

Appendix B: Angular integral

The integrals in \textbf{B1} and \textbf{B4} are solved in this appendix, whose forms are

$$I_1 = \int_{0}^{\pi} d\theta \frac{\sin^2 \theta}{(a - b \cos \theta)}; \quad \text{(B1)}$$

$$I_2 = \int_{0}^{\pi} d\theta \frac{\sin^2 \theta}{(a - b \cos \theta)^2},$$

where

$$a = p^2 + k^2, \quad b = 2pk. \quad \text{(B2)}$$

They both have the same structure, namely

$$I_n = \int_{0}^{\pi} d\theta \frac{\sin^2 \theta}{(a - b \cos \theta)^n}, \quad \text{(B3)}$$

and are evaluated in the same way.

First write it as an integral from 0 to 2$\pi$ so that it can be converted into a contour integral along a closed circular path in the complex plane. Since $\cos \theta = \cos(2\pi - \theta)$ and $\sin^2 \theta = \sin^2(2\pi - \theta)$ then $I_n$ can equally be written as

$$I_n = \int_{0}^{2\pi} d\theta \frac{\sin^2(2\pi - \theta)}{(a - b \cos(2\pi - \theta))^n}$$

$$= \int_{\pi}^{2\pi} d\xi \frac{\sin^2 \xi}{(a - b \cos \xi)^n}, \quad \text{(B4)}$$

such that

$$I_n = \frac{1}{2} \int_{0}^{2\pi} d\theta \frac{\sin^2 \theta}{(a - b \cos \theta)^n}. \quad \text{(B5)}$$

Now write this as a closed integral over the variable $z = e^{i\theta}$ in the complex plane, with $\sin \theta = (-i/2)(z + z^{-1})$ and $\cos \theta = (1/2)(z + z^{-1})$:

$$I_n = \frac{i}{8} \oint dz \frac{2^n z^{n-3} (z^2 - 1)^2}{(2az - bz^2 - b)^n}. \quad \text{(B6)}$$

Write the denominator of the integrand as

$$(2az - bz^2 - b)^n = (-b)^n (z - z_1)^n (z - z_2)^n, \quad \text{(B7)}$$

where

$$z_1 = \frac{a + \sqrt{a^2 - b^2}}{b}, \quad z_2 = \frac{a - \sqrt{a^2 - b^2}}{b}, \quad \text{(B8)}$$

to bring the integral into the form

$$I_n = 2^{n-3} i \oint dz \frac{z^{n-3} (z^2 - 1)^2}{(-b)^n (z - z_1)^n (z - z_2)^n}. \quad \text{(B9)}$$

Now the integral has the familiar form

$$\oint dz \frac{f(z)}{(z - z_1)^m (z - z_2)^m}$$

where the contour is the unit circle with $f(z)$ analytic and continuous at $z_1$ and $z_2$, and it can be solved by the standard method of summing over the residues of $f(z)$.

Note carefully the location of the poles. From \textbf{B2} it follows that

$$a - b = p^2 + k^2 - 2pk = (p - k)^2 > 0, \quad \text{(B10)}$$

and there it is always true that

$$a > b. \quad \text{(B11)}$$

Accordingly

$$z_1 = \frac{k}{p} \theta(k - p) + \frac{p}{k} \theta(p - k), \quad \text{(B12)}$$

$$z_2 = \frac{p}{k} \theta(k - p) + \frac{k}{p} \theta(p - k). \quad \text{(B13)}$$

In conclusion $|z_2| < 1$ while $|z_1| > 1$, so the pole at $z_1$ is discounted because it is not inside the unit circle. This leaves two poles in the expression \textbf{B9}: one at $z = 0$ and one at $z = z_2$. 
For $n = 2$ the integral in \[B9\], by the Cauchy formula, has the form

\[
I_2 = \frac{i}{2} \frac{(2\pi i)^4}{b^2} \left[ \left. \frac{1}{(z^2-1)^4} \right|_{z=0} + \frac{d}{dz} \left. \frac{(z^2-1)^3}{(z^2-1)^4} \right|_{z=2} \right] \\
= -\frac{\pi}{(2pk)^2} \left[ \left. \frac{1}{z_1^2} \frac{(z_2-1)^2}{z_2(z_2-1)} \right|_{z=0} - \frac{2}{z_2} \right] + \frac{4}{(2z_2-1)^2} , \tag{B14}
\]

while for $n = 1$ it is

\[
I_1 = \frac{i}{4} \frac{(2\pi i)^4}{(b)} \left[ \left. \frac{d}{dz} \frac{(z^2-1)^2}{(z^2-1)^3} \right|_{z=0} + \frac{(z^2-1)^3}{z_2(z_2-1)} \right|_{z=2} \right] \\
= \frac{\pi}{4pk} \left[ \frac{z_1 + z_2}{z_1^2 z_2^2} + \frac{(z_2^2-1)}{z_2^2(z_2-1)} \right] , \tag{B15}
\]

Finally substitute the explicit forms of $z_1$ and $z_2$ to obtain

\[
I_1 = \frac{\pi}{2k^2} \theta(k-p) + \frac{\pi}{2p^2} \theta(p-k) , \tag{B16}
\]

and

\[
I_2 = -\frac{\pi}{2k^2(p^2-k^2)} \theta(k-p) - \frac{\pi}{2p^2(k^2-p^2)} \theta(p-k) . \tag{B17}
\]

\[\]

[1] Latham Boyle and Neil Turok. “Cancelling the vacuum energy and Weyl anomaly in the standard model with dimension-zero scalar fields”. [2110.06258.pdf] 2021.

[2] Sergei Alexandrov. Immirzi parameter and fermions with non-minimal coupling. Class. Quant. Grav., 25:145012, 2008.

[3] Peter W. Higgs. Broken symmetries, massless particles and gauge fields. Phys. Lett., 12:132–133, 1964.

[4] Peter W. Higgs. Spontaneous Symmetry Breakdown without Massless Bosons. Phys. Rev., 145:1156–1163, 1966.

[5] T. W. B. Kibble. Symmetry breaking in non-Abelian gauge theories. Phys. Rev., 155:1554–1561, 1967.

[6] F. Englert and R. Brout. Broken Symmetry and the Mass of Gauge Vector Mesons. Phys. Rev. Lett., 13:321–323, 1964.

[7] G. S. Guralnik, C. R. Hagen, and T. W. B. Kibble. Global Conservation Laws and Massless Particles. Phys. Rev. Lett., 13:585–587, 1964.

[8] Steven Weinberg. Precise relations between the spectra of vector and axial vector mesons. Phys. Rev. Lett., 18:507–509, 1967.

[9] S. L. Glashow. Partial Symmetries of Weak Interactions. Nucl. Phys., 22:579–588, 1961.

[10] Steven Weinberg. A Model of Leptons. Phys. Rev. Lett., 19:1264–1266, 1967.

[11] A. Salam. “Relativistic Groups and Analyticity”. In 8th Nobel Symposium on Elementary Particle Theory, 1968, page 367, 1968.

[12] H. David Politzer. Reliable Perturbative Results for Strong Interactions? Phys. Rev. Lett., 30:1346–1349, 1973.

[13] R. N. Cahn and G. Goldhaber. The experimental foundations of particle physics. Cambridge Univ. Press, Cambridge, 2009.

[14] H. Fritzsch, Murray Gell-Mann, and H. Leutwyler. Advantages of the Color Octet Gluon Picture. Phys. Lett. B, 47:365–368, 1973.

[15] M. P. Nakada, J. D. Anderson, C. C. Gardner, J. McClure’, and C. Wong. Neutron Spectrum from p+d Reaction. Phys. Rev., 110:594–595, 1958.

[16] Yoichiro Nambu. Quasiparticles and Gauge Invariance in the Theory of Superconductivity. Phys. Rev., 117:648–663, 1960.

[17] Julian S. Schwinger. Gauge Invariance and Mass. Phys. Rev., 125:397–398, 1962.

[18] Serguei Chatrchyan et al. Search for the Standard Model Higgs Boson Produced in Association with $W$ and $Z$ Bosons in $pp$ Collisions at $\sqrt{s} = 7$ TeV. JHEP, 11:088, 2012.

[19] Chiara Mariotti. Observation of a new Boson at a Mass of 125 gev With the CMS Experiment at the LHC. In 13th Marcel Grossmann Meeting on Recent Developments in Theoretical and Experimental General Relativity, Astrophysics, and Relativistic Field Theories, pages 352–372, 2015.

[20] Serguei Chatrchyan et al. Observation of a New Boson at a Mass of 125 GeV with the CMS Experiment at the LHC. Phys. Lett. B, 716:30–61, 2012.
[21] Serguei Chatrchyan et al. A Search for a Doubly-Charged Higgs Boson in $pp$ Collisions at $\sqrt{s} = 7$ TeV. Eur. Phys. J. C, 72:2189, 2012.

[22] Georges Aad et al. Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC. Phys. Lett. B, 716:1–29, 2012.

[23] Eleni Mountricha. Search for the Standard Model Higgs boson in the $H \rightarrow ZZ(*) \rightarrow 4l$ channel with the ATLAS experiment at the LHC leading to the observation of a new particle compatible with the Higgs boson. PhD thesis, Natl. Tech. U., Athens and Orsay, 10 2012.

[24] Kenneth Lane. “Two Lectures on Technicolor”. [arxiv.org/abs/hep-ph/0202225] 2002.

[25] H. Terazawa. t quark mass predicted from a sum rule for lepton and quark masses. Phys. Rev. D, 22:2921–2921, 1980. [Erratum: Phys.Rev.D 41, 3541 (1990)].

[26] H. Terazawa, Y. Chikashige, K. Akama, and T. Matsuki. Simple Relation Between the Fine Structure and Gravitational Constants. Phys. Rev. D, 15:1181, 1977.

[27] William A. Bardeen, Christopher T. Hill, and Manfred Lindner. Minimal Dynamical Symmetry Breaking of the Standard Model. Phys. Rev. D, 41:1647, 1990.

[28] W. J. Marciano. HEAVY TOP QUARK MASS PREDICTIONS. Phys. Rev. Lett., 62:2793–2796, 1989.

[29] Christopher T. Hill. Topcolor: Top quark condensation in a gauge extension of the standard model. Phys. lett. B, 266:419–424, 1991.

[30] Christopher T. Hill and Elizabeth H. Simmons. Strong Dynamics and Electroweak Symmetry Breaking. Phys. Rept., 381:235–402, 2003. [Erratum: Phys.Rept. 390, 553–554 (2004)].

[31] V. A. Miransky, Masaharu Tanabashi, and Koichi Yamawaki. Is the t Quark Responsible for the Mass of W and Z Bosons? Mod. Phys. Lett. A, 4:1043, 1989.

[32] R. Sekhar Chivukula, Bogdan A. Dobrescu, Howard Georgi, and Christopher T. Hill. Top Quark Seesaw Theory of Electroweak Symmetry Breaking. Phys. Rev. D, 59:075003, 1999.

[33] Bogdan A. Dobrescu and Christopher T. Hill. Electroweak symmetry breaking via top condensation seesaw. Phys. Rev. Lett., 81:2634–2637, 1998.

[34] Hsin-Chia Cheng, Bogdan A. Dobrescu, and Jiayin Gu. Higgs Mass from Compositeness at a Multi-Tev Scale. JHEP, 08:095, 2014.

[35] Leonard Susskind. Dynamics of Spontaneous Symmetry Breaking in the Weinberg-Salam Theory. Phys. Rev. D, 20:2619–2625, 1979.

[36] Steven Weinberg. Implications of Dynamical Symmetry Breaking. Phys. Rev. D, 13:974–996, 1976. [Addendum: Phys.Rev.D 19, 1277–1280 (1979)].

[37] Savas Dimopoulos and Leonard Susskind. Mass Without Scalars. Nucl. Phys. B, 155:237–252, 1979.

[38] Estia Eichten and Kenneth D. Lane. Dynamical Breaking of Weak Interaction Symmetries. Phys. Lett. B, 90:125–130, 1980.

[39] Thomas Appelquist, Anuradha Ratnaweera, John Terning, and L. C. R. Wijewardhana. The Phase structure of an SU(N) gauge theory with N(f) flavors. Phys. Rev. D, 58:105017, 1998.

[40] Sven Bjarke Gudnason, Chris Kouvaris, and Francesco Sannino. Towards working technicolor: Effective theories and dark matter. Phys. Rev. D, 73:115003, 2006.

[41] Kenneth D. Lane. An Introduction to technicolor. In Theoretical Advanced Study Institute (TASI 93) in Elementary Particle Physics: The Building Blocks of Creation - From Microfermions to Megaparsecs, 6 1993.

[42] Michael J. Dugan, Howard Georgi, and David B. Kaplan. Anatomy of a Composite Higgs Model. Nucl. Phys. B, 254:299–326, 1985.

[43] Maxim Perelstein. Little Higgs models and their phenomenology. Prog. Part. Nucl. Phys., 58:247–291, 2007.

[44] N. Arkani-Hamed, A. G. Cohen, E. Katz, and A. E. Nelson. The Littlest Higgs. JHEP, 07:034, 2002.

[45] Ian Low, Witold Skiba, and David Tucker-Smith. Little Higgses from an antisymmetric condensate. Phys. Rev. D, 66:072001, 2002.

[46] C. F. Berger, M. Perelstein, and F. Petriello. Top quark properties in little Higgs models. In 2005 International Linear Collider Physics and Detector Workshop and 2nd ILC Accelerator Workshop, 12 2005.

[47] Michele Redi and Andrea Tesi. Implications of a Light Higgs in Composite Models. JHEP, 10:166, 2012.

[48] Abdesslam Arhib, Rachid Benbrik, Hicham Harouiz, Stefano Moretti, Yan Wang, and Qi-Shu Yan. Implications of a light charged Higgs boson at the LHC run III in the 2HDM. Phys. Rev. D, 102(11):115040, 2020.

[49] David B. Kaplan. Flavor at SSC energies: A New mechanism for dynamically generated fermion masses. Nucl. Phys. B, 365:259–278, 1991.

[50] Michele Redi and Andreas Weiler. Flavor and CP Invariant Composite Higgs Models. JHEP, 11:108, 2011.

[51] Deog Ki Hong and Du Hwan Kim. Composite (pseudo) scalar contributions to $m_\nu \approx 2$. Phys. Lett. B, 758:370–372, 2016.

[52] Shinya Matsuzaki and Koichi Yamawaki. Walking from 750 GeV to 950 GeV in the technipion zoo. Phys. Rev. D, 93(11):115027, 2016.

[53] Keisuke Harigaya and Yasunori Nomura. A Composite Model for the 750 GeV Diphoton Excess. JHEP, 03:091, 2016.

[54] Yuichiro Nakai, Ryosuke Sato, and Kohsaku Tobioka. Footprints of New Strong Dynamics via Anomaly and the 750 GeV Diphoton. Phys. Rev. Lett., 116(15):151802, 2016.

[55] Roberto Franceschini, Gian F. Giudice, Jernej F. Kamnik, Matthew McCullough, Alex Pomarol, Riccardo Rattazzi, Michele Redi, Francesco Riva, Alessandro Strumia, and Riccardo Torre. What is the $\gamma\gamma$ resonance at 750 GeV? JHEP, 03:144, 2016.

[56] Emiliano Molinaro, Francesco Sannino, and Natascia Viguaro. Minimal Composite Dynamics versus Axion Origin of the Diphoton excess. Mod. Phys. Lett. A, 31(26):1650155, 2016.

[57] Ligong Bian, Ning Chen, Da Liu, and Jing Shu. Hidden confining world on the 750 GeV diphoton excess. Phys. Rev. D, 93(9):095011, 2016.

[58] Yang Bai, Joshua Berger, and Ran Lu. 750 GeV dark pion: Cousin of a dark G -parity odd WIMP. Phys. Rev. D, 93(7):076009, 2016.

[59] James M. Cline and Zuowei Liu. LHC diphotons from electroweakly pair-produced composite pseudoscalars. 12 2015.

[60] P. Ko and Takaaki Nomura. Dark sector shining through 750 GeV dark Higgs boson at the LHC. Phys. Lett. B, 758:205–211, 2016.

[61] Michele Redi, Alessandro Strumia, Andrea Tesi, and Elena Vigiana. Di-photon resonance and Dark Matter as heavy pions. JHEP, 05:078, 2016.
Keisuke Harigaya and Yasunori Nomura. Hidden Pion Varieties in Composite Models for Diphoton Resonances. Phys. Rev. D, 94(7):075004, 2016.

Robert Foot and John Gargalionis. Explaining the 750 GeV diphoton excess with a colored scalar charged under a new confining gauge interaction. Phys. Rev. D, 94(1):011703, 2016.

Sho Iwamoto, Gabriel Lee, Yael Shadmi, and Robert Ziegler. Diphoton Signals from Colorless Hidden Quarkonia. Phys. Rev. D, 94(1):015003, 2016.

Yang Bai, Joshua Berger, James Osborne, and Ben A. Stefanek. Phenomenology of Strongly Coupled Chiral Gauge Theories. JHEP, 11:153, 2016.

Archil Kobakhidze, Fei Wang, Lei Wu, Jin Min Yang, and Mengchao Zhang. 750 GeV diphoton resonance in a top and bottom seesaw model. Phys. Lett. B, 757:92–96, 2016.

Bob Holdom. Raising Condensates Beyond the Ladder. Phys. Lett. B, 213:365–369, 1988.

B. Holdom. HOLDOM REPLIES TO: COMMENT ON ‘CONTINUUM LIMIT OF QUENCHED THEORIES.’. Phys. Rev. Lett., 63:1889, 1989.

Koichi Yamawaki. Dynamical symmetry breaking with large anomalous dimension. In 14th Symposium on Theoretical Physics: Dynamical Symmetry Breaking and Effective Field Theory, 3 1996.

Jarno Rantaharju, Claudio Pica, and Francesco Sannino. Conformal Window of Gauge Theories with Four-Fermion Interactions and Ideal Walking. Phys. Rev. D, 82:035021, 2010.

Jarno Rantaharju, Claudio Pica, and Francesco Sannino. Ideal Walking Dynamics via a Gauged NJL Model. Phys. Rev. D, 96(1):014512, 2017.

Roshan Foadi, Mads T. Frandsen, and Francesco Sannino. 125 GeV Higgs boson from a not so light technicolor scalar. Phys. Rev. D, 87(9):095001, 2013.

Francesco Sannino and Kimmo Tuominen. Orientifold theory dynamics and symmetry breaking. Phys. Rev. D, 71:051901, 2005.

Dennis D. Dietrich and Francesco Sannino. Conformal window of SU(N) gauge theories with fermions in higher dimensional representations. Phys. Rev. D, 76:055018, 2007.

Dennis D. Dietrich, Francesco Sannino, and Kimmo Tuominen. Light composite Higgs from higher representations versus electroweak precision measurements: Predictions for CERN LHC. Phys. Rev. D, 72:055001, 2005.

Giacomo Cacciapaglia, Claudio Pica, and Francesco Sannino. Fundamental Composite Dynamics: A Review. Phys. Rept., 877:1–70, 2020.

A. Palatini. Deduzione invariantiva delle equazioni gravitazionali dal principio di Hamilton. Rend. Circ. Mat. Palermo, 43:203–212.

Carlo Rovelli and Francesca Vidotto. Covariant Loop Quantum Gravity: An Elementary Introduction to Quantum Gravity and Spinfoam Theory. Cambridge Monographs on Mathematical Physics. Cambridge University Press, 11 2014. See Eq.(3.30) p.64.

Soren Holst. Barbero’s Hamiltonian derived from a generalized Hilbert-Palatini action. Phys. Rev. D, 53:5966–5969, 1996.

Alejandro Perez and Carlo Rovelli. Physical effects of the Immirzi parameter. Phys. Rev. D, 73:044013, 2006.

Carlo Rovelli and Francesca Vidotto. Covariant Loop Quantum Gravity: An Elementary Introduction to Quantum Gravity and Spinfoam Theory. Cambridge Monographs on Mathematical Physics. Cambridge University Press, 11 2014. See Eq.(3.32)-(3.34) p.65.

M. A. Zubkov. Torsion instead of Technicolor. Mod. Phys. Lett. A, 25:2885–2898, 2010.

A. S. Belyaev and I. L. Shapiro. Torsion action and its possible observables. Nucl. Phys. B, 543:20–46, 1999.

I. L. Shapiro. The Interaction of quantized matter fields with torsion and metric: NJL model. Mod. Phys. Lett. A, 9:729–733, 1994.

I. L. Shapiro. Physical aspects of the space-time torsion. Phys. Rept., 357:113, 2002.

Carlo Rovelli and Francesca Vidotto. Covariant Loop Quantum Gravity: An Elementary Introduction to Quantum Gravity and Spinfoam Theory. Cambridge Monographs on Mathematical Physics. Cambridge University Press, 11 2014. See Eq.(3.20) p.63.

Carlo Rovelli. “Quantum gravity“. Cambridge Monographs on Mathematical Physics. Univ. Pr., Cambridge, UK, 2004. See Eq.(2.33) p.38.

M. A. Zubkov. Schwinger–Dyson equation and NJL approximation in massive gauge theory with fermions. Annals Phys., 354:72–88, 2015.

V. A. Miransky. Dynamical symmetry breaking in quantum field theories. 1994. See §8.4, specifically Eq.(8.49) therein.