A UNIFIED CONFORMAL MODEL 
FOR FUNDAMENTAL INTERACTIONS 
WITHOUT DYNAMICAL HIGGS FIELD 

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Abstract 

A Higgsless model for strong, electro–weak and gravitational interactions is proposed. This model is based on the local symmetry group $SU(3) \times SU(2)_L \times U(1) \times C$ where $C$ is the local conformal symmetry group. The natural minimal conformally invariant form of total lagrangian is postulated. It contains all Standard Model fields and gravitational interaction. Using the unitary gauge and the conformal scale fixing conditions we can eliminate all four real components of the Higgs doublet in this model. However the masses of vector mesons, leptons and quarks are automatically generated and are given by the same formulas as in the conventional Standard Model. In this manner one gets the mass generation without the mechanism of spontaneous symmetry breaking and without the remaining real dynamical Higgs field. The gravitational sector is analyzed and it is shown that the model admits in the classical limit the Einsteinian form of gravitational interactions.

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1 Introduction

The recent evidence for top quark production with the top mass estimated as $m_t = 174 \pm 10^{+12}_{-12} \text{GeV}$ \cite{1} implies that the Higgs particle – if exists – may have the mass of the order of 1TeV: in fact the central value of $m_H$ implied by the present data of $m_t$ and $m_{W}$ was estimated by Hioki and Najima \cite{2} at $m_H \approx 1700 \text{GeV}$ with an enormous error however. Since in the lowest order $\lambda = \frac{1}{2}(m_H v^2)$ one can afraid that the Higgs self-coupling $\lambda$ would be also very large ($\lambda \approx 25$ for the central value of $m_H$ given by Hioki and Najima). Such strong Higgs self-interaction would mean that the loops with Higgs particles would dominate all other contributions. Therefore the perturbative predictions in Standard Model(SM) for many quantities become unreliable. Hence the predictive power of the SM and its consistency may be questionable.

The Higgs particle with such a large mass becomes suspicious. It is natural therefore to search for a modification of SM in which all confirmed by experiment particles would exist but the Higgs particle as the observed object would be absent.

We show in this work that such a modification of SM is possible under the condition that one joints to strong and electro–weak interactions also the gravitational interaction. This extension of the class of SM interactions is in fact very natural. Indeed whenever we have the strong and electro–weak interactions of elementary particles, nuclea, atoms or other objects we have also at the same time the gravitational interactions. It seems natural therefore to consider an unified model for strong, electro–weak and gravitational interactions which would describe simultaneously all four fundamental interactions. It is well known that gravitational interactions give a negligible effect to most of strong or electro–weak elementary particle processes. We show however that they may play the crucial role in a determination of the physical fields and their masses in the unified model and that their presence allows to eliminate all Higgs fields from the final lagrangian.

In turn we recall that in the conventional Standard Model the Higgs mechanism of spontaneous symmetry breaking (SSB) provides a simple and effective instrument for mass generation of weak gauge bosons, quarks and leptons. However, despite of many efforts of several groups of experimentalists \cite{3} the postulated Higgs particle of the SM was not observed. Hence one might expect that the model for strong and electro–weak interactions
supplemented by the gravitational interaction in which all dynamical Higgs fields may be eliminated can provide a natural framework for a description of elementary particle fundamental interactions.

In order to construct a new form of total lagrangian for the theory of strong and electro–weak interactions extended by the gravitational interactions we observe that the gauge symmetry $SU(3) \times SU(2)_L \times U(1)$ of the fundamental interactions may be naturally extended by the local conformal symmetry. The choice of the unitary gauge condition for $SU(2)_L$ gauge group allows to eliminate the three out of four Higgs fields from the complex Higgs doublet. In turn the choice of the scale fixing condition connected with the local conformal symmetry allows to eliminate the last Higgs field. In that manner all four Higgs fields can be gauged away completely! It is remarkable that in spite of the elimination of all Higgs fields in our model the vector meson, lepton and quark masses are generated and at the tree level they are given by the same analytical formulas as in the conventional SM.

Thus it may be that the dynamical real Higgs field and the associated Higgs particles are in fact absent and it is therefore not surprising that they could not be detected in various experiments.

We review in Section 2 the present problems with a very massive Higgs particle. Next in Section 3 we discuss the properties of local conformal symmetry and its representations in field space of arbitrary spin. We present in Section 4 the form of the total lagrangian of our unified theory of electro–weak, strong and gravitational interactions determined by the gauge and the local conformal invariance. The noteworthy feature of the obtained lagrangian is the lack of the Higgs mass term $\mu^2 \Phi^\dagger \Phi$. We show next that using the unitary gauge condition and the conformal scale fixing condition we can eliminate all dynamical Higgs fields from the theory! We show in Section 5 that in spite of the lack of dynamical Higgs fields the masses of vector mesons, leptons and quarks are generated and at the tree level are given by the same analytical expressions in terms of coupling constants as in the conventional SM.

We stress that the renormalizability of our model depends on the value of the new coupling constant $\beta$ which determines the properties of gravitational sector. We discuss in Section 6 the variant of our model with $\beta \neq 0$. This leads to the model with massive vector mesons which is nonrenormalizable. In order to get definite perturbative predictions – especially for electro–weak processes – we have to introduce the ultraviolet cutoff $\Lambda$. We show the
close connection between the large Higgs mass $m_H$ and $\Lambda$. We illustrate this relation in the case of universal electro–weak parameters $\varepsilon_{N1}$, $\varepsilon_{N2}$ and $\varepsilon_{N3}$ of Altarelli et al. [4] for which we show that the difference between SM results for $\varepsilon_{N1}$ and in our model is essentially proportional to $\log \frac{\Lambda^2}{m_H^2}$; thus if one chooses $\Lambda \approx m_H$ one obtains the same analytical formulas for $\varepsilon_{N1}$ in SM and in our Higgsless model. We show also how using so called General Equivalence Theorems one can calculate the high energy limit for various processes in our model.

We present in Section 7 the analysis of the gravitational sector in the unified model. We show that our unified model after determination of the unitary gauge and scale fixing leads already at classical level to the conventional gravitational theory with Einstein–Hilbert lagrangian implied by the conformal Penrose term contained in the unified lagrangian.

We present in Section 8 the special version of our model (with $\beta = 0$) which may lead to perturbatively renormalizable model of fundamental interactions. We discuss shortly some open problems of this formulation of the unified theory.

Finally we discuss in Section 9 three alternatives for a description of fundamental interactions which are given by the conventional SM or its extensions, Higgsless renormalizable SM and nonrenormalizable Higgsless models. We discuss also some open problems connected with derivation of predictions in low and high energy regions from nonrenormalizable Higgsless models.

The present work is the extension of our previous paper [5] and contains the answer to several questions raised by its readers.

\section{Difficulties with Standard Model Higgs particle.}

We shall argue that the recently announced evidence for the top quark with the mass

$$m_t = 174 \pm 10^{+13}_{-12} GeV$$

may lead to a serious conceptual and calculational problems in the Standard Model. The relatively heavy top quark with the mass (2.1) – heavier than expected on the base of LEP1-CDF-UA1 data [1], [3] – shifts up the expected
region of SM Higgs mass and consequently also the area of expected Higgs quartic self-coupling $\lambda$. The analysis of the value of Higgs mass following from the one-loop formula for the $W$–meson mass carried out by Hioki and Najima \[2\] leads to the central value

\[ m_H \approx 1700 GeV. \quad (2.2) \]

Since the Higgs self–coupling constant $\lambda$ and the Higgs mass are connected at the tree level by the formula

\[ \lambda = \frac{1}{2} \left( \frac{m_H}{\langle \phi \rangle} \right)^2, \quad \langle \phi \rangle = 246 GeV \quad (2.3) \]

the value (2.2) implies that

\[ \lambda \approx 25. \quad (2.4) \]

This looks very dangerous; however to be honest we should mention that within the present experimental errors for $m_t$ given by FNAL result and for other experimental quantities being the input for the estimation (2.2) there is a considerable admissible deviation for $m_H$ from the central value 1700 GeV \[7\] \[8\]. Consequently $m_H$ and therefore also $\lambda$ may be much smaller.

Despite the fact that the present electro-weak data are not very conclusive the result (2.1) compels many authors to consider the possibilities of large Higgs mass and strong Higgs self–coupling more seriously. The super–strong Higgs self–coupling (like (2.4) or even smaller) would evidently break–down the perturbative calculations for many processes for which Higgs loops with $\lambda$-coupling contributes. For instance the two-loop perturbation expansion for the partial width decay $\Gamma(H \rightarrow \bar{f}f)$ of the Higgs particle into the fermion – antifermion pair can be written in the form

\[ \Gamma(H \rightarrow \bar{f}f) = \Gamma_0[1 + 0.11\left(\frac{m_H}{1 TeV}\right)^2 - 0.78\left(\frac{m_H}{1 TeV}\right)^4] \quad (2.5) \]

where $\Gamma_0$ is the partial width in the Born approximation and the second and third term in the bracket represent the one- and the two-loop contributions respectively \[10\].

We see that with increasing $m_H$ the importance of the two-loop contribution rapidly increases: for $m_H > 375 GeV$ the two-loop contribution
dominates the one-loop and for $m_H > 1200\text{GeV}$ the width becomes negative! This demonstrates the complete breakdown of perturbation theory for the Higgs mass of the order of 1TeV.

We see therefore that the supposition that the real Higgs field and the corresponding Higgs particle exists in the SM may lead to rather fundamental conceptual and calculational difficulties. Therefore it seems justified at present to look for a modification of SM in which all experimentally confirmed facts would be reproduced but the Higgs particle as the observed object would not exist.

Recently there were proposed several Higgsless models for electro–weak and strong interactions. In particular Shildknecht and collaborators proposed the Higgsless massive vector boson model [11] and they have compared some of its predictions with the predictions of the conventional SM. In the work [12] it was proposed a Higgsless SM with nonrenormalizable current–current and dipol–dipol interactions. Finally in [13] it was proposed a gauged $\sigma$–model for electro–weak interactions.

It seems to us that our Higgsless model based on the extension of electro–weak and strong interactions by gravitational interactions which leads to the extension of gauge symmetry by the local conformal symmetry presents a most natural frame–work for fundamental interactions.

3 Local conformal symmetry

Let $M^{3,1}$ be the pseudo–Riemannian space time with the metric $g_{\alpha\beta}$ with the signature $(+,−,−,−)$. Let $\Omega(x)$ be a smooth strictly positive function on $M^{3,1}$. Then the conformal transformation in $M^{3,1}$ is defined as the transformation which changes the metric by the formula

$$g_{\mu\nu}(x) \rightarrow \tilde{g}_{\mu\nu}(x) = \Omega^2(x)g_{\mu\nu}(x).$$

(3.1)

The set of all conformal transformations forms the multiplicative abelian infinite–dimensional group $C$ with the obvious group multiplication law.

It is evident from (3.1) that $(M^{3,1}, g_{\mu\nu})$ and $(M^{3,1}, \tilde{g}_{\mu\nu})$ have identical causal structure and conversely it is easy to show that any two space times which have identical causal structure must be related by a local conformal transformation.
The conformal transformations occur in many problems in general relativity. In particular Canuto et. al. proposed the scale–covariant theory of gravitation, which provides an interesting alternative for the conventional Einstein theory [14].

It should be stressed that a conformal transformation is not a diffeomorphism of space time. The physical meaning of the conformal transformations follows from the transformation law of the length element

\[
dl(x) = \sqrt{-g_{ij}dx^i dx^j} \rightarrow \tilde{dl}(x) = \Omega(x) dl(x).
\]

Hence a local conformal transformation changes locally the length scale. Since in some places of the Earth one utilizes the meter as the length scale, whereas in other places one utilizes the feet or the ell as the length scale one may say that one utilizes the local conformal transformations in everyday live. Similarly one verifies that the conformal transformation changes locally the proper time

\[
ds(x) = \sqrt{g_{\mu\nu}dx^\mu dx^\nu} \rightarrow \tilde{ds}(x) = \Omega(x) ds(x).
\]

Since the physical phenomena should be independent of the unit chosen locally for the length, the proper time, mass etc. the group \( C \) of local conformal transformations should be a symmetry group of physical laws.

In order to avoid any confusion we stress that the group \( C \) has nothing in common with the 15 parameter conformal group \( SO(4,2) \) defined locally in the \( M^{3,1} \) by the action of Poincare, dilatation and special conformal transformations.

Comparing the physical meaning of local conformal transformations and the local gauge \( SU(2)_L \) transformations of SM associated with the concept of the weak isospin it seems that the conformal transformations are not less natural symmetry transformation than the nonabelian gauge transformations in the SM.

We shall give now a construction of the representation of the conformal group \( C \) in the field space. Let \( \Psi \) be a tensor or spinor field of arbitrary spin. Define the map

\[
\Omega \rightarrow U(\Omega)
\]

by the formula

\[
\bar{\Psi}(x) = U(\Omega)\Psi(x) = \Omega^s(x)\Psi(x), \quad s \in \mathbb{R}
\]
The number $s$ is determined by the condition of conformal invariance of field equation. We say that field equation for $\Psi$ is conformal invariant if there exist $s \in \mathbb{R}$ such that $\Psi(x)$ is a solution with the metric $g_{\mu\nu}(x)$ if and only if $\tilde{\Psi}(x)$ given by (3.3) is a solution with the metric $\tilde{g}_{\mu\nu}(x)$. The number $s$ is called the conformal weight of $\Psi$ \cite{13}, \cite{16}, \cite{17}. It is evident that the map $\Omega \to U(\Omega)$ defines the representation of $C$ in the field space.

Using the above definitions one can calculate the conformal weight for a field of arbitrary spin. Let for instance $F_{\mu\nu}$ be the Maxwell field on $(M^{3,1}, g)$ which satisfies the equation

$$g^{\mu\sigma} \nabla_\sigma F_{\mu\nu} = 0$$
$$\nabla_{[\sigma} F_{\mu\nu]} = 0.$$  

Using the definition of the covariant derivative $\tilde{\nabla}_\sigma$ with respect to $\tilde{g}_{\mu\nu}$ metric and (3.3) one obtains

$$\tilde{g}^{\mu\sigma} \tilde{\nabla}_\sigma (\Omega^s F_{\mu\nu}) = (n - 4 + s)\Omega^{s-3} g^{\mu\sigma} F_{\mu\nu} \nabla_\sigma \Omega$$
$$\tilde{\nabla}_{[\sigma} (\Omega^s F_{\mu\nu]}) = s\Omega^{s-1} (\nabla_{[\sigma} \Omega) F_{\mu\nu]}.$$  

We see that for $n \neq 4$ the Maxwell equations are not conformally invariant. For $n = 4$ the Maxwell equations are invariant if the conformal weight $s$ equals to zero.

Similarly one can show that the Yang–Mills field strength $F_{\mu\nu}^a$ has the conformal weight $s = 0$ whereas the massless Dirac field has the conformal weight $s = -\frac{3}{2}$. It is noteworthy that the scalar massless field $\Phi$ satisfying the Laplace–Beltrami equation

$$\Delta \Phi = 0$$

is not conformal invariant. In fact it was discovered by Penrose that one has to add to the Lagrangian on $(M^{3,1}, g)$ the term

$$-\frac{1}{6} R \Phi^2$$

where $R$ is the Ricci scalar, in order that the corresponding field equation is conformal invariant with the conformal weight $s = -1$ \cite{18}.  

7
4 A unified model for strong, electro–weak and gravitational interactions

We postulate that the searched unified theory of strong, electro–weak and gravitational interactions will be determined by the condition of invariance with respect to the group $G$

$$G = SU(3) \times SU(2)_L \times U(1) \times C$$  \hspace{1cm} (4.1)

where $C$ is the local conformal group defined by (3.1). Let $\Psi$ be the collection of vector meson, fermion and scalar fields which appear in the conventional minimal SM for electro–weak and strong interactions. Then the minimal natural conformal and $SU(3) \times SU(2)_L \times U(1)$–gauge invariant total lagrangian $L(\Psi)$ may be postulated in the form:

$$L = [L_G + L_F + L_Y + L_{\Phi} + \beta \partial_\mu |\Phi| \partial^\mu |\Phi| - \frac{1}{6}(1 + \beta) R \Phi^\dagger \Phi + L_{grav}] \sqrt{-g}$$  \hspace{1cm} (4.2)

Here $L_G$ is the total lagrangian for the gauge fields $A_a^\mu$, $W^b_\mu$ and $B_\mu$, $a = 1, \ldots, 8$, $b = 1, 2, 3$ associated with $SU(3) \times SU(2)_L \times U(1)$ gauge group

$$L_G = -\frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} - \frac{1}{4} W_{\mu\nu}^b W^{b\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu},$$  \hspace{1cm} (4.3)

and $F_{\mu\nu}^a$, $W_{\mu\nu}^b$ and $B_{\mu\nu}$ are the conventional field strengths of gauge fields in which the ordinary derivatives are replaced by the covariant derivatives e.g.

$$B_{\mu\nu} = \nabla_\mu B_\nu - \nabla_\nu B_\mu,$$ \hspace{1cm} (4.4)

e tc.; $L_F$ is the lagrangian for fermion field interacting with the gauge fields; $L_Y$ represents the Yukawa interactions of fermion and scalar fields; $L_{\Phi}$ is the lagrangian for the scalar fields

$$L_{\Phi} = (D\Phi)^\dagger (D\Phi) - \lambda (\Phi^\dagger \Phi)^2$$ \hspace{1cm} (4.5)

where $D$ denotes the covariant derivative with connections of all symmetry groups. Notice that the condition of conformal invariance does not admit the Higgs mass term $\mu^2 \phi^\dagger \phi$ which assures the mechanism of spontaneous symmetry breaking and mass generation in the conventional formulation.
The term

\[ \beta \partial_\mu |\Phi| \partial^\mu |\Phi| \]  

(4.6)
is gauge invariant. It may be surprising that (4.6) depends on |\Phi|. Observe however that the lagrangian \( L_\phi \) can be written in the form

\[ (D\Phi)^\dagger (D\Phi) = \partial_\mu |\Phi| \partial^\mu |\Phi| + |\Phi|^2 L_\sigma (g(\Phi), W, B) \]  

(4.7)where \( L_\sigma (g(\Phi), W, B) \) is a gauged–sigma–model–like lagrangian and

\[ \Phi = \begin{pmatrix} \phi_u \\ \phi_d \end{pmatrix} = g(\Phi) \begin{pmatrix} 0 \\ |\Phi| \end{pmatrix}, \quad g(\Phi) = \frac{1}{|\Phi|} \begin{pmatrix} \bar{\phi}_d & \phi_u \\ -\bar{\phi}_d & \phi_u \end{pmatrix} \]  

(4.8)where \( g(\Phi) \) is \( SU(2)_L \) gauge unitary matrix.

We see therefore that the term like (4.6) is already present in the conventional \( L_\Phi \) lagrangian.

The term

\[ -\frac{1}{6} (1 + \beta) R \Phi^\dagger \Phi \]  

(4.9)
is the Penrose term, which assures that the lagrangian (4.2) is conformal invariant.

The last term in (4.2) is the Weyl term

\[ L_{\text{grav}} = -\rho C^2, \quad \rho > 0, \]  

(4.10)where \( C^\delta_{\alpha\beta\gamma} \) is the Weyl tensor which is conformally invariant. Using the Gauss–Bonnet identity we can write \( C^2 \) in the form

\[ C^2 = 2(R^\mu\nu R_{\mu\nu} - \frac{1}{3} R^2). \]  

(4.11)

We see that the condition of conformal invariance does not admit in (4.2) the conventional gravitational Einstein lagrangian

\[ L = \kappa^{-2} R \sqrt{-g}, \quad \kappa^2 = 16\pi G. \]  

(4.12)

It was shown however by Stelle \[19\] that quantum gravity sector contained in (4.2) is perturbatively renormalizable whereas the quantum gravity defined by the Einstein lagrangian (4.12) coupled with matter is nonrenormalizable \[20\]. Hence, for a time being it is an open question which form of gravitational
interaction is more proper on the quantum level. We show in Section 7 that the Einstein lagrangian (4.12) may be reproduced by Penrose term if the physical scale is properly determined. In Section 8 we discuss the role of quantum effects which may reproduce the lagrangian (4.12) and give the classical Einstein theory as the effective induced gravity.

Notice that conformal symmetry implies that all coupling constants in the present model are dimensionless.

The theory given by (4.2) is our conformally invariant proposition alternative to the standard Higgs–like theory with SSB. Its new, most important feature is the local conformal invariance. It means that simultaneous rescaling of all fields (including the field of metric tensor) with a common, arbitrary, space–time dependent factor $\Omega(x)$ taken with a proper power for each field (the conformal weight) will leave the Lagrangian (4.2) unaffected. The symmetry has a clear and obvious physical meaning [21], [18]. It changes in every point of the space–time all dimensional quantities (lengths, masses, energy levels, etc) leaving theirs ratios unchanged. It reflects the deep truth of the nature that nothing except the numbers has an independent physical meaning.

The freedom of choice of the length scale is nothing but the gauge fixing freedom connected with the conformal symmetry group. In the conventional approach we define the length scale in such a way that elementary particle masses are the same for all times and in all places. This will be the case when we rescale all fields with the $x$–dependent conformal factor $\Omega(x)$ in such a manner that the length of the rescaled scalar field doublet is fixed i.e.

$$\tilde{\Phi}^+ \tilde{\Phi} = \frac{v^2}{2} = \text{const.} \quad (4.13)$$

(We shall discuss the problem of mass generation in details in Section 5.)

The scale fixing for the conformal group (4.13) is distinguished by nothing but our convenience. Obviously we can choose other gauge fixing condition, e.g. we can use the freedom of conformal factor to set

$$\sqrt{-\tilde{g}} = 1; \quad (4.14)$$

this will lead to other local scales but it will leave physical predictions unchanged.
Consider, for example the scale fixing condition (4.14). Imposing (4.14) on the conformal invariant theory given by (4.2) we obtain the lagrangian \( \tilde{L}(\tilde{\Psi}, \tilde{V}_\mu, \tilde{\Phi}, \tilde{g}_{\mu\nu}) \) describing dynamics of the fields \( \tilde{\Psi}, \tilde{V}_\mu, \tilde{\Phi}, \tilde{g}_{\mu\nu} \). The arguments of \( \tilde{L} \) stand for all fermion, vector, scalar and tensor fields of the model and fulfill the condition (4.14). \( \tilde{L} \) is no longer conformal invariant as the scale was fixed by (4.14). We can change variables of \( \tilde{L} \) according to the rule

\[
\tilde{\Psi} = \left( \frac{\sqrt{2} |\tilde{\Phi}|}{v} \right)^{-3/2} \tilde{\Psi} \quad (4.15a)
\]

\[
\tilde{V}_\mu = \tilde{V}_\mu \quad (4.15b)
\]

\[
\tilde{g}_{\mu\nu} = \left( \frac{\sqrt{2} |\tilde{\Phi}|}{v} \right)^2 \tilde{g}_{\mu\nu} \quad (4.15c)
\]

\[
\tilde{\Phi} = \left( \frac{\sqrt{2} |\tilde{\Phi}|}{v} \right)^{-1} \tilde{\Phi} \quad (4.15d)
\]

where \( \tilde{g} \) is no longer restricted but \( \tilde{\Phi} \) fulfills (4.13) what follows from (4.15d). Such a change of variable is an example of conformal transformation but, as was said, it is not a symmetry of \( \tilde{L} \). In fact we have

\[
\tilde{L}(\tilde{\Psi}(\tilde{\Psi}, \tilde{V}_\mu, \tilde{\Phi}, \tilde{g}_{\mu\nu}), \tilde{V}_\mu(\tilde{\Psi}, \tilde{V}_\mu, \tilde{\Phi}, \tilde{g}_{\mu\nu}), \tilde{\Phi}(\tilde{\Psi}, \tilde{V}_\mu, \tilde{\Phi}, \tilde{g}_{\mu\nu}), \tilde{g}_{\mu\nu}(\tilde{\Psi}, \tilde{V}_\mu, \tilde{\Phi}, \tilde{g}_{\mu\nu})) =
\]

\[
= \tilde{L}(\tilde{\Psi}, \tilde{V}_\mu, \tilde{\Phi}, \tilde{g}_{\mu\nu}) \quad (4.16)
\]

where \( \tilde{L}(\tilde{\Psi}, \tilde{V}_\mu, \tilde{\Phi}, \tilde{g}_{\mu\nu}) \) is the lagrangian which one would obtain by imposing the scale fixing condition (4.13) directly on (4.2). It should be stressed that the functional form of \( \tilde{L} \) in terms of its arguments is different than \( \tilde{L} \) of its arguments (compare (5.1) and (8.2) for concrete examples). In such a sense theories obtained from different scale fixings are mathematically equivalent. They will be equivalent also physically if identifications of physical and mathematical objects in the theories being compared will be consistent with theirs mathematical equivalence. For example if we assume that \( \tilde{g} \) describes physical metric we cannot assume that this metric is described also by \( \tilde{g} \).
5 Generation of lepton, quark and vector boson masses

We demonstrate now that using the conformal group scale fixing condition (4.13) we can generate the same lepton, quark and vector meson masses as in the conventional SM without however use of any kind of Higgs mechanism and SSB.

In fact inserting the scale fixing condition (4.13) into the Lagrangian (4.2) we obtain

\[
\tilde{L} = L^{scaled} = [L_G + L_F + L_\phi^{scaled} + L_Y^{scaled} - \frac{1}{12}v^2 R + L_{grav}]\sqrt{-g},
\]

in which the condition (4.13) was inserted into \(L_\phi\) and \(L_Y\). We should use the symbol \(\tilde{\Phi}, \tilde{\Psi}\) etc. for the rescaled fields in (5.1), however for the sake of simplicity we shall omit “~” sign over fields in the following considerations.

The condition (4.13) together with the unitary gauge fixing of \(SU(2)_L \times U(1)\) gauge group, reduce by (4.8) the Higgs doublet to the form

\[
\Phi^{gauge} = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}, \quad v > 0
\]

and produce the tree level mass terms for leptons, quarks and vector bosons associated with \(SU(2)_L\) gauge group. For instance the \(\Phi–\)lepton Yukawa interaction \(L_Y^l\)

\[
L_Y^l = - \sum_{i=e,\mu,\tau} G_i \bar{l}_i R (\Phi^* l_i L) + h.c.
\]

passes into

\[
L_Y^{gauged} = - \frac{1}{\sqrt{2}} v (G_e \bar{e} e + G_\mu \bar{\mu} \mu + G_\tau \bar{\tau} \tau)
\]

giving the conventional, space–time independent lepton masses

\[
m_e = \frac{1}{\sqrt{2}} G_e v, \quad m_\mu = \frac{1}{\sqrt{2}} G_\mu v, \quad m_\tau = \frac{1}{\sqrt{2}} G_\tau v.
\]

Similarly one generates from \(\Phi–\)quark Yukawa interaction \(L_Y^q\) the corresponding quark masses. In turn from \(L_\phi\)-lagrangian (4.5) using the gauge condition (5.2) one obtains

\[
(D_\mu \tilde{\Phi})^\dagger D^\mu \tilde{\Phi} = \frac{g_2^2 v^2}{4} W_\mu^+ W^{-\mu} + \frac{g_1^2 + g_2^2}{8} v^2 Z^2
\]
where

\[ Z_\mu = -\sin \theta_W B_\mu + \cos \theta_W W^3_\mu, \quad \cos \theta_W = \frac{g_2}{\sqrt{g_1^2 + g_2^2}}. \]

Hence one obtains the following vector mesons masses

\[ m_W = \frac{v}{2}g_2, \quad m_Z = \frac{m_W}{\cos \theta_W}. \] (5.5)

It is remarkable that the analytical form for tree level fermion and vector meson masses in terms of coupling constants and the parameter \( v \) is the same as in the conventional SM. We see therefore that the Higgs mechanism and SSB is not indispensable for the fermion and vector mesons mass generation!

We note that the fermion–vector boson interactions in our model are the same as in SM. Hence analogously as in the case of conventional formulation of SM one can deduce the tree level relation between \( v \) and \( G_F \) – the four–fermion coupling constant of \( \beta \)-decay:

\[ v^2 = (2G_F)^{-1} \rightarrow v = 246 \text{GeV}. \] (5.6)

Here we have used the standard decomposition \( g^{\mu\nu} \sqrt{-g} = \eta^{\mu\nu} + \kappa' h^{\mu\nu} \) (see e.g. [22]) which reduces the tree level problem for the matter fields to the ordinary flat case task.

We see therefore that the resulting expressions for masses of physical particles are identical as in the conventional SM.

Let us stress that the scale fixing condition like (4.13) does not break \( SU(2)_L \times U(1) \) gauge symmetry. The symmetry is broken (or rather one of gauge equivalent description is fixed) when (4.13) is combined with unitary gauge condition of electro–weak group leading to (5.2). However, also after imposing of a gauge condition like (5.2) we have a remnant of both the conformal and \( SU(3) \times SU(2)_L \times U(1) \) initial gauge symmetries: this is reflected in the special, unique relations between couplings and masses in our model.
6 Precision tests of electro–weak interactions and high energy behavior in the present model.

Our model represents in fact the gauge field theory model with massive vector mesons and fermions. It is well-known that such models are in general nonrenormalizable [23]. We remind however that in nonrenormalizable Fermi model for weak interactions we can make a definite predictions for low energy phenomena e.g. for $\mu$ or neutron decays. Similarly the recent progress with so called Generalized Equivalence Theorem allows to make definite predictions for the scattering operator in nonrenormalizable models like gauged nonlinear $\sigma$–model or other nonrenormalizable gauge field theory models [24]. Hence in our model we can obtain definite predictions for electro–weak phenomena if we consider processes with energy $\sqrt{s}$ below some ultraviolet (UV) cutoff $\Lambda$. We wish to demonstrate that the cutoff $\Lambda$ is determined by the Higgs mass $m_H$ appearing in the Standard Model. Hence, from this point of view, Higgs mass is nothing else as the UV cutoff which assures that the truncated perturbation series is meaningful. We shall try to elucidate this problem on the example of so called precision tests of electro–weak theory.

One–loop radiative corrections to various electro–weak quantities or processes can be expressed in terms of three quantities $\Delta r$, $\Delta \rho$ and $\Delta k'$. We refer to the recent excellent review by Kniehl for the precise definitions of these quantities and for their analytical expressions [25]. For an illustration we recall that the expression for $W$–meson mass, up to one loop order, has the form

$$M_W = \frac{M_Z}{\sqrt{2}} \left\{ 1 + \sqrt{1 + \frac{2\sqrt{2}\pi \alpha}{M_Z G_F (1 - \Delta r)}} \right\}^\frac{1}{2}$$

where

$$\alpha = \frac{1}{137.036},$$

$$G_F = 1.16639 \times 10^{-5} GeV^{-2},$$

$$M_Z = 91.1899 \pm 0.0044 GeV$$

and $\Delta r(m_t, m_H)$ is the one loop correction to $\mu$–decay amplitude which in Standard Model depends on top and Higgs masses. Taking the experimental
value for W–mass \( M_W = 80.21 \pm 0.18 \text{GeV} \) and the recently reported top mass \( m_t = 174 \pm 17 \text{GeV} \) one gets from (6.1) the central value of Higgs mass \( m_H \approx 1700 \text{GeV} \) with the error of several hundreds of GeV [2].

It was suggested by Altarelli et. al [4] to pass from \( \Delta r, \Delta \rho \) and \( \Delta k' \) to new quantities \( \varepsilon_{N1}, \varepsilon_{N2} \) and \( \varepsilon_{N3} \) such that \( \varepsilon_{N2} \) and \( \varepsilon_{N3} \) depend on \( m_t \) only logarithmically. These parameters characterize the degree of \( SU(2)_L \times U(1) \) symmetry breaking and their numerical value significantly different from zero would signal a ”new physics” [11][4].

If we calculate these parameters in our model in one–loop approximation we find the specific class of Feynman diagrams with fermion and vector boson loops which contributes to them. Since some vector boson loops will produce divergences, e.g. in the case of fermion – massive vector boson coupling constant, one has to introduce either the new renormalization constant or UV cutoff \( \Lambda \) which can be given by the formula [11]

\[
\log \frac{\Lambda^2}{\mu^2} = \frac{2}{4-D} - \gamma_E + \log 4\pi
\]

where \( \mu \) is the reference mass of dimensional regularization, \( D \) is the space–time dimension and \( \gamma_E \) is the Euler’s constant.

One obtains the formula for \( \varepsilon_{Ni} \) parameters in SM if one adds to the class of Feynman diagrams in our model all appropriate one–loop diagrams with Higgs internal lines. Using the results of [11] and [26] one obtains

\[
\varepsilon^{SM}_{N1} - \varepsilon^{CSM}_{N1} = \frac{3\alpha(M_Z^2)}{16\pi c_0^2} \log \left( \frac{\Lambda^2}{m_H^2} \right) + \mathcal{O}\left( \frac{M_Z^2}{m_H^2} \log \left( \frac{M_H^2}{M_Z^2} \right) \right)
\]

\[
\varepsilon^{SM}_{N2} - \varepsilon^{CSM}_{N2} = \mathcal{O}\left( \frac{M_Z^2}{m_H^2} \log \left( \frac{M_H^2}{M_Z^2} \right) \right)
\]

\[
\varepsilon^{SM}_{N3} - \varepsilon^{CSM}_{N3} = \frac{\alpha(M_Z^2)}{48\pi s_0^2} \log \left( \frac{\Lambda^2}{m_H^2} \right) + \mathcal{O}\left( \frac{M_Z^2}{m_H^2} \log \left( \frac{M_H^2}{M_Z^2} \right) \right)
\]

where \( CSM \) index of \( \varepsilon_{Ni} \) means that the quantity was calculated in our Conformal Standard Model. Here \( \alpha(M_Z^2) = \frac{1}{129} \) and \( c_0 \) and \( s_0 \) are defined by the formula

\[
\frac{s_0^2(1-s_0^2)}{s_0^2 c_0^2} = \frac{\pi \alpha(M_Z^2)}{\sqrt{2} G_F M_Z^2}
\]

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The above formulas indicate a role which plays in SM the very large Higgs mass: first the term \(O\left(\frac{M_Z^2}{m_H^2} \log \left(\frac{M_H^2}{M_Z^2}\right)\right)\) for \(m_H > 1\) TeV can be disregarded and second if we take the UV cutoff \(\Lambda \simeq m_H\) then the prediction for \(\varepsilon_{Ni}\)–parameters in the conventional SM and our nonrenormalizable model coincide. Thus the very large Higgs mass preferred by the top mass \(m_t = 174\) GeV plays in the conventional SM the role of UV cutoff parameter. If the Higgs particle will be not found then our model provides an extremely natural frame–work for the description of electro–weak and strong interactions at least up to TeV energies.

We would like to discuss now the problem of getting predictions from our nonrenormalizable model for electro–weak and strong interactions considered in the flat space–time. Take the process \(A + B \rightarrow C + D\) in our model. This process – up to L–loop order – will be described by the corresponding Feynman diagrams with A, B, C and D external lines and some number of internal fermion, massive vector mesons, gluon and photon lines. Since theory is nonrenormalizable one has to introduce the proper UV cut off \(\Lambda\).

The problem of elaboration of an effective calculational scheme for our model is considerably facilitated by the fact that introducing the suitable Stueckelberger auxiliary fields we can transform our model into the gauged nonlinear \(\sigma\)–model (GNL\(\sigma\)M) (see e.g. [11], [13] and the discussion in Section 9). It is known that perturbative calculations in GNL\(\sigma\)M with cutoff \(\Lambda\) are well elaborated and lead to interesting physical predictions for various processes [11], [24].

In fact it was recently shown that so called General Equivalence Theorem (GET) holds in gauge field theories irrespectively if they are renormalizable or nonrenormalizable [24]. This remarkable theorem can be applied in the case of SM for heavy Higgs at high energy where

\[m_H, E \gg M_W, m_{f_i}\]

where \(E\) is the total energy and \(m_{f_i}\) are lepton and quark masses respectively. It was shown that the leading parts coming from the L–loop diagrams are those diagrams for which \(N\) defined as

\[N = \text{power of } m_H + \text{power of } E\]

becomes maximal. Using GET one relatively easily determines the leading contribution for any L–loop in SM and obtains high energy limit of a given
scattering amplitude [24]. In the case of Higgsless nonrenormalizable gauge field theory model one introduces cutoff Λ: in this case at high energy limit defined by inequalities

$$\Lambda > E \gg M_W, m_f,$$

the leading diagrams are those for which

$$N = \text{power of } \Lambda + \text{power of } E$$

is maximal. Comparing (6.5) with (6.6) we see as in the case of the $\varepsilon_{N_1}$-parameters that the UV cutoff Λ in Higgsless gauge models replaces the large mass $m_H$. Using the criterion (6.6) and GET one obtains the high energy limit of scattering amplitude for various processes also in the non-renormalizable gauge models, like e.g. in the Higgsless GNLσM [24].

We see therefore that nonrenormalizability does not prevent us from getting definite predictions for physical processes in the low or high energy region from our model. Consequently the nonrenormalizable Higgsless models may be as a useful in description of experimental data as the conventional SM.

We considered hence the general variant of our model with $\beta \neq 0$ which leads to nonrenormalizable gauge field theory. However the special case of our model with $\beta = 0$ discussed in Section 8 gives a renormalizable model for fundamental interactions.

## 7 Gravity Sector

Let us impose the scale fixing condition (4.13) on the lagrangian (4.2) and collect all gravitational terms. The lagrangian reads:

$$L^{scaled} = [L^{scaled}_{matter} - \frac{1}{12}(1 + \beta)v^2R - 2\rho(R^{\mu\nu}R_{\mu\nu} - \frac{1}{3}R^2) - \frac{\lambda}{4}v^4]\sqrt{-g}$$

where we have selected the part $L^{scaled}_{matter}$ (describing the matter interacting with gravity) from the remaining purely gravitational terms.

The variation of (7.1) with respect to the metric $g^{\mu\nu}$ leads to the following classical equation of motion:

$$\rho[-\frac{2}{3}R_{;\mu;\nu} + 2R_{\mu\nu;\eta} - \frac{2}{3}g_{\mu\nu}R^{\eta;\eta;\eta} -$$

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\[4R^\eta\lambda R_{\mu\eta\nu\lambda} + \frac{4}{3}RR_{\mu\nu} + g_{\mu\nu}(R^\eta\lambda R_{\eta\lambda} - \frac{1}{3}R^2)] + \]
\[
\frac{1}{12}(1 + \beta)v^2(R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R) + \frac{\lambda}{8}v^4g_{\mu\nu} = \frac{1}{2}T_{\mu\nu}. \tag{7.2}
\]

In the empty case \(T_{\mu\nu} = 0\) this equation is satisfied by all solutions of an empty space Einstein equation with a properly chosen cosmological constant \(\Lambda\):

\[R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R + \Lambda g_{\mu\nu} = 0. \tag{7.3}\]

In fact (7.3) implies that

\[R_{\mu\nu} \sim g_{\mu\nu} \quad \Rightarrow \quad R_{\mu\nu} = \frac{1}{4}Rg_{\mu\nu} \tag{7.4}\]

and then

\[R_{\mu\nu} = \Lambda g_{\mu\nu}. \tag{7.5}\]

Inserting (7.4) into (7.2) we find that the part proportional to \(\rho\) vanishes. The remnant can be collected leading to the relation

\[
\frac{1}{8}v^2g_{\mu\nu}(\frac{2}{3}(1 + \beta)\Lambda - \lambda v^2) = 0 \tag{7.6}
\]

where the empty space condition \(T_{\mu\nu} = 0\) were used for the right hand side of (7.6).

It is easy to conclude that (7.6) is satisfied when

\[
\Lambda = \frac{3}{2(1 + \beta)}\lambda v^2. \tag{7.7}
\]

Equation (7.7) relates \(\lambda\) with a potentially observable cosmological constant \(\Lambda\).

Let us go back to the case with the matter. Observe that the term linear in the curvature appears in (7.1) with the coefficient \(-\frac{1}{12}(1 + \beta)v^2\). In the case of \(\beta = 0\) we have the old–fashion Standard Model minimally conformally coupled with gravity. In this case, in comparison with the Newtonian constant entering to the ordinary Einstein’s theory (4.12) the coefficient \(-\frac{1}{12}v^2\) standing in front of \(R\) in (7.1) has an opposite sign and is smaller of many orders of magnitude \((v^2\kappa^2 \approx 10^{-38})\). If it would be the only purely gravitational
term in the theory it will mean that the geometry in tree approximation is
generated by the negative energy and with an extremal strength. This is the
price we would have to pay for the positive kinetic term of scalar fields in
(4.5), for the gauge invariance and for the renormalizability of the matter
sector.

If we want to reproduce the correct gravitational sector already at the
classical level rather than preserve renormalizability of the material sector
we have to admit for nonzero $\beta$ coupling. This would lead us to a model
which is equivalent to the nonrenormalizable gauged nonlinear sigma model
in the material sector. Accepting this price we can put

$$-\frac{1}{12}(1 + \beta)v^2 = \kappa^{-2}$$

reproducing the Newtonian coupling in front of curvature $R$ in (7.1). This
would mean that $\beta \approx -10^{38}$. Notice however that taking the scale fixing
condition (4.13) the term $\beta \partial_{\mu} |\Phi| \partial^\mu |\Phi|$ vanishes. Hence it looks like that the
only role of this term is to generate the proper value of Newton constant in
the Einstein–Hilbert tree level lagrangian resulting from the Penrose term.
We will go back to this point in Section 8.

8 Towards the renormalizable theory.

We have shown in Section 6 that the nonrenormalizability of our model does
not rise serious calculational problems within the energy range presently ac-
cessible in experiments. What more, if Higgs particle will be not found then
it cannot be generally excluded that nonrenormalizability will be the indisp-
ensable feature of every realistic particle theory model. It would mean that
we are compelled to work with a theory valid for a limited energy regions and
even for limited classes of phenomena. Clearly this situation is unsatisfactory
and people will always try to find a general description scheme unifying dif-
ferent phenomena and independent on the considered energy range. Hopes
for such a universal description are usually set on renormalizable unified
models. Being motivated by these hopes let us go back to the problem of
renormalizability of our model.

It is easy to see that nonrenormalizability of the matter part of lagrangian
(4.2) is connected with the presence of nonlinear interaction (4.6). To see
this we can approximate (4.2) demanding that
\[ g^{\mu\nu} = \eta^{\mu\nu}. \]  

(8.1)

This is a conformally flat approximation rather than the flat approximation as we have the scale fixing freedom, and the part of relations (8.1) can be understood as making use of this freedom e.g. in the form of condition (4.14). Putting (4.14) into (4.2) we obtain
\[ \tilde{L} = L_{\text{unimodular}} = L_G + L_F + L_Y + L_\Phi + \beta \partial_\mu |\Phi| \partial^{\mu} |\Phi| - \frac{1}{6} (1 + \beta) R \Phi^\dagger \Phi + L_{\text{grav}} \]

(8.2)

Putting in turn (8.1) we obtain the conformally flat approximation lagrangian \( L_{\text{cfa}} \):
\[ L_{\text{cfa}} = L_G + L_F + L_Y + L_\Phi + \beta \partial_\mu |\Phi| \partial^{\mu} |\Phi| \]

(8.3)

For \( \beta = 0 \) this is just the renormalizable SM lagrangian (without the negative scalar mass term \( -\mu^2 \Phi^\dagger \Phi \) however).

The presence of \( \beta \)-term was justified in Section 7 by the condition of the proper Einsteinian limit of the theory at its classical level. This led us to the rather large value \( |\beta| \sim \frac{m_{\text{PLANCK}}^2}{2^4} \). Fortunately considerations of Section 7 showed that the predictions in particle sector of our theory are insensitive to this huge value of \( \beta \)-coupling.

As we have mention in Section 7 for the case with \( \beta = 0 \) after the choice of physical scaling the obtained tree level Newton constant is not correct. It was suggested by various authors that the corrected Newton constant in quantum theories of gravity may be obtained by inclusion of radiative corrections [27] [28] (see [29] for a pedagogical introduction and [30] for the recent review of the subject). The authors of [31] have discussed a wide class of theories which contains also our model in the case of \( \beta = 0 \). They have shown that taking the proper values for the nonobserved coupling constants like \( \rho \) or \( \lambda \) and renormalization scale one may generate the induced Newton and cosmological constants with experimental values. However this method is – in our opinion – incomplete since the problem of mass values of elementary particles in the framework in which gravitational constants were determined was not considered. In fact the value of the Newton constant has not an absolute meaning. This constant disappear from the empty space Einstein
equations. In the presence of matter the value of Newton constant can be rescaled with simultaneous rescaling of masses and energy levels. Thus the value of induced Newton and cosmological constants must be compared with the effective masses of classical matter fields obtained within the same level of perturbative analysis before going to the final conclusions on induced Einstein lagrangian. According to our knowledge the quantum gravity corrected expressions for the effective masses were not derived so far in the literature.

Until this problem is solved one cannot conclude that that the renormalizable model for fundamental interactions with $\beta = 0$ is physically meaningful.

### 9 Discussion.

The elementary particle physics is at present at a crossroad. We have in fact three drastically different alternatives:

I$^0$ The Higgs particle exists, its mass will be experimentally determined and will have the value predicted by the radiative corrections of SM. This will confirm the SSB mechanism for mass generation, the validity of SM framework and it will represent an extraordinary success of quantum SM framework.

II$^0$ The Higgs particle exists but its mass is considerably different from that predicted by the radiative corrections of SM. This would signal some kind of ”New Physics” which will imply a reformulation of the present version of SM.

III$^0$ The Higgs particle does not exist. This will lead to a rejection of SM with Higgs sector and it will give preference to Higgsless models for fundamental interactions. In this situation we have two general possibilities:

IIIA$^0$ The physical Higgsless model is renormalizable. The example of such model was discussed by us in Section 8.

IIIB$^0$ The physical Higgsless model is nonrenormalizable. It may be that the renormalizability of Quantum Gravity determined by Einstein–Hilbert action integral coupled with matter fields is not an ”accident at work in quantum field theory” but it represents a universal feature that physical fundamental interactions considered simultaneously are nonrenormalizable. In this situation we are compelled to use the nonrenormalizable models of quantum field theory for a description of fundamental interactions and we have to learn how to deduce predictions for experiments from such models.
Several nonrenormalizable models for electroweak interactions were proposed like Schildknecht et al. model, [1], GNLσM [13], or gauge field theory models with condensates [32]. We have presented in Section 4 a new unified nonrenormalizable model for fundamental interactions based on the gauge and local conformal symmetry.

Our model – in spite of its nonrenormalizability – provides the definite predictions for low and very high energy interactions in terms of the parameters of the model, energy $E$ and the cutoff $\Lambda$. The direct calculations of electro–weak parameters $\varepsilon_{N1}$, $\varepsilon_{N2}$ and $\varepsilon_{N3}$ demonstrate that the Standard Model and the present model results differ by the term proportional to $\log \frac{\Lambda^2}{m_H^2}$: thus it looks like that the very high Higgs mass $m_H$ plays in SM the role of the UV cutoff which in the present model may be replaced by parameter $\Lambda$. We see therefore that the predictive power of our model may be comparable with that of the conventional SM.

In view of the possibility that nonrenormalizable nonabelian massive gauge field theories have to be used for a description of fundamental interactions it seems necessary to develop perturbative and nonperturbative methods for extracting predictions for scattering amplitudes and observables from such models. In particular one should develop the corresponding Generalized Equivalence Theorems and determine explicitly the high energy behavior of cross sections in such models. The comparison of the obtained results with analytic formulas coming from Lipatov calculations [33] would be very inspiring. It would be also useful to develop systematic two–loop calculus with UV cutoff $\Lambda$ for electro–weak processes. We plan in a near future to present several examples of such calculations.

The present model allows to obtain the Einsteinian form of gravitational interactions in the classical limit. It can be also analyzed by means of effective action for induced gravity [30].

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