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UNIFIED COMPOSITENESS
OF LEPTONS, QUARKS AND HIGGS BOSONS

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Abstract

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The unified compositeness of leptons, quarks and Higgs bosons is proposed as a possible
scenario for New Physics beyond the Standard Model. The following topics of the scenario are
briefly discussed:
- Chiral gauge exceptional symmetry E_6 as a strong internal binding mechanism;
- Higgs doublet as a composite Goldstone boson;
- Nonlinear Standard Model as a prototype "low energy" effective field theory of the unified
compositeness;
- Hidden local symmetry and an improved "low energy" effective field theory of the unified
compositeness;
- Heavy composite vector bosons and vector boson dominance of the SM gauge interactions;
- Universal dominant residual interactions as a signature of the unified compositeness;
- Manifestations of the residual interactions and the potential of the future TeV colliders to
uncover the unified compositeness.

Аннотация

Кабаченко В.В., Пирогов Ю.Ф. Объединенная композитность лептонов, кварков и хиггсов-
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В качестве возможного сценария новой физики за пределами стандартной модели пред-
лагается объединенная композитность лептонов, кварков и хиггсовских бозонов. Кратко
обсуждаются следующие разделы этого сценария:
- Киральная калибровочная исключительная симметрия E_o как механизм сильной связи;
- Хиггсовский дублет как составной голдстоуновский бозон;
- Нелинейная стандартная модель как прототип "низкоэнергетической" эффективной тео-
рии объединенной композитности;
- Скрытая локальная симметрия и улучшенная "низкоэнергетическая" эффективная теория
объединенной композитности;
- Тяжелые составные векторные бозоны и векторно-бозонная доминантность калибровоч-
ных взаимодействий СМ;
- Универсальные доминантные остаточные взаимодействия как сигнатура объединенной
композитности;
- Проявления остаточных взаимодействий и возможность обнаружения объединенной ком-
позитности на будущих e^+e^- линейных коллайдерах тэвных энергий.

Introduction

Are leptons and quarks composite or not? This is the question. The same is for Higgs bosons. If both these types of the Standard Model (SM) fields were composite simultaneously, having common substructure, one could, in principle, solve a lot of the SM problems. First of all, considering leptons and quarks as light composite fermions one could find a rationale for the well-known generation problem and that of the fermion quantum numbers (see, e.g., [1]). Further, treating the SM Higgs doublet as composite Goldstone boson [2] one could solve the naturalness problem [3] of the SM without supersymmetry. More than that, one could also try to unify Higgs self-interactions and Yukawa interactions reducing thus their arbitrariness. This interactions should emerge as residual ones from an interplay of the hyperstrong binding gauge interactions and the perturbative weak gauge interactions, etc.

In a series of papers [4]–[6] one of the present authors (Yu.F.P.) has developed the scheme of the unified compositeness of leptons, quarks and Higgs bosons (gauge bosons being still elementary) as a promising scenario for New Physics beyond the SM. Some phenomenological consequences of the scenario have been further studied by us in a number of subsequent papers [7,8]. In this report we present a brief survey of these developments.

1. Scenario of unified compositeness

Let us expose the principal ideas of the scenario of unified compositeness of leptons, quarks and Higgs bosons. The scenario encounters two stages: that of the dynamical symmetry breaking and that of the spontaneous symmetry breaking.

Dynamical symmetry breaking. Let a hypothetical hyperstrong gauge theory S_{loc}, responsible for the tight internal binding of the SM composite particles, possess a global chiral symmetry G, that of the Lagrangian. As a result of the nontrivial topological structure in S_{loc} at a scale $\mathcal{F}$ a partial dynamical symmetry breaking $G \to H$ takes place, where H is a residual symmetry of the vacuum $|0>$. It is supposed that H embeds the symmetry of the SM: $H \supseteq I_{SM} = SU(2)_L \times U(1)_Y$. In this, the broken symmetry

G/H corresponds to true Goldstone bosons, in particular, to the Higgs doublet. On the other hand, the unbroken chiral symmetry H is responsible via the 't Hooft anomaly matching condition for the appearance of the massless composite fermions in addition to the massless composite Higgs bosons.

In reality, part of the Lagrangian symmetry G is gauge: $G \supset I_{loc} \supset I_{loc}^{SM}$. Thus this symmetry undergoes the partial dynamical braking as $I_{loc} \to R_{loc}$ with the residual gauge symmetry being $R_{loc} = I_{loc} \cap H$ and $R_{loc} \supseteq I_{loc}^{SM}$. The Goldstone bosons corresponding to the broken part of the local symmetry I_{loc}/R_{loc} are absorbed via the Higgs mechanism by the proper gauge bosons V, the latter becoming massive: $M_V \simeq g_V \mathcal{F}$. The rest of the Goldstone bosons, including the Higgs doublet, is still massless at this stage.

In the framework of the effective field theory the dynamical symmetry breaking is described by the nonlinear model G/H [9,10] with intrinsic gauge theory I_{loc} being spontaneously broken as $I_{loc} \to R_{loc}$. Here the weak gauge interactions are considered as a perturbation unimportance in the lowest approximation for the basic properties of the symmetry breaking pattern. Still, interactions of I_{loc} explicitly violate symmetry G, and their account results in important physical effects.

Spontaneous symmetry breaking. At I_{loc} being turned off, all the possible orientations of the residual symmetry H inside the total symmetry G are equivalent. This results in a set of the degenerate H invariant vacua $|\xi> \equiv \xi |0>$, where $\xi \in G/H$. Explicit violations due to I_{loc} being turned on, this equivalence is lost. The question arises as to which the preferred orientation of H relative to I_{loc} is and what the true vacuum of the lowest energy is. This is the so-called vacuum alignment problem [11,12]. The orientation in question is determined by the radiative corrections caused by the virtual emission and absorption of the I_{loc} gauge bosons.

Namely, let us consider the effective action $\Gamma(\phi) = -V(\phi)_{eff} + \ldots$, where the one-loop effective potential is $V(\phi)_{eff} = \mu^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2 + \mathcal{O}((\phi^\dagger \phi)^3)$. The following general statement is true [11,12]. The radiative corrections due to the dynamically unbroken symmetry R_{loc} contribute positively to μ^2 and thus stabilize the unperturbed vacuum $|0>$ trying to orient the residual global symmetry H along the symmetry R_{loc} itself. And v.v., the radiative corrections due to dynamically broken symmetry I_{loc}/R_{loc} contribute negatively to μ^2 and hence destabilize vacuum $|0>$ trying to disorient symmetry H relative to symmetry R_{loc}. The net effect is $\mu^2 = C\bar{g}^2/(4\pi)^2 \mathcal{F}^2$, where $\bar{g}$ is a generic effective gauge coupling constant, $\mathcal{F}$ is the mass scale of the dynamical symmetry breaking and C is a numerical being determined by the ratio of the two contributions with opposite signs.

Thus there are tree possibilities for the curvature μ^2 of the effective potential.

- $\boldsymbol{\mu^2 \geq 0}$. This is the case of the convex potential. Here R_{loc} is left unbroken, and Goldstone bosons ϕ turn into the pseudo-Goldstone ones with $m^2 \geq 0$.
- $\boldsymbol{\mu^2 < 0}$. In this case the potential is concave. This means that the symmetry R_{loc} is spontaneously broken. Goldstone bosons ϕ turn into the would-be Goldstone ones and are absorbed through the Higgs mechanism by the corresponding weak gauge bosons.
- $\boldsymbol{\mu^2 = 0}$. This is degenerate case of the flat one-loop potential. Here one-loop approximation is insufficient and two-loop corrections have to be considered [13].

We take for granted that one of the two last cases is realized, and the spontaneous symmetry breaking of R_{loc} takes place. The Higgs self-interactions described by $V_{eff} \simeq V_H$ are no longer fundamental but arise as residual ones from more fundamental gauge interactions. Similarly, the effective Yukawa interactions must arise as a result of the radiative effects due to I_{loc}. Both these types of interactions are deprived of their fundamental status. Thus a kind of the unification of the Higgs and Yukawa interactions occurs, and there appears an opportunity, at least, in principle, to reduce their arbitrariness.

Now, there are two possibilities for the mass scale $\mathcal{F}$ of the unified compositeness.

- **TeV compositeness.** It realizes in the most general one-loop case. Here one can show that $\mathcal{F} = \mathcal{O}(v)$, where v is the SM v.e.v. We consider this as phenomenologically unacceptable. For $\mathcal{F} \gg v$ to take place a fine tuning is required, and this is unnatural.

- **Deca-TeV compositeness.** For the theory to be natural, one should put to one-loop $v \equiv 0$. Then in two-loops one has $\mu^2 = \mathcal{O}((g^2/(4\pi)^2)^2 \mathcal{F}^2)$, whereas $\lambda = \mathcal{O}(g^2/(4\pi)^2)$ as before. It follows hereof that $v \equiv \mu/\sqrt{\lambda} = \mathcal{O}(g/4\pi\,\mathcal{F})$, or $v = \mathcal{O}(\sqrt{\alpha_W/4\pi}\mathcal{F})$. In other words, one has $\mathcal{F} = \mathcal{O}(m_W/\alpha_W)$, or $\mathcal{F} = \mathcal{O}(10\,\text{TeV})$. Thus, the natural two-loop hierarchy $\mathcal{F} \gg v$ between compositeness scale and Fermi scale arises.

It is this last scenario that will be developed in what follows. More details can be found in refs. [4]–[6].

2. Chiral gauge exceptional symmetry

A paramount problem in building a realistic composite model of leptons and quarks is to find underlying forces capable of binding these particles at the distances much smaller than their Compton wave lengths. Strongly coupled non-Abelian gauge theories S_{loc} provide presently a unique well-fitted framework for such a binding mechanism.

It is imperative that in the process of confinement a set of (almost) massless composite fermions should emerge. In other words, this is to require that some residual chiral symmetry should be left unbroken in the transition. A necessary (but not sufficient) condition for this is the chiral anomaly matching condition [3,14].

There are conclusive arguments that strongly interacting $SU(N)$ gauge theories with n Dirac constituent fermions (so that their $SU(N)$ representation is vector-like) break the chiral symmetry $SU(n) \times SU(n) \times U(1)$ down to the vector-like one $SU(n) \times U(1)$ and hence do not produce massless composite fermions [15,16]. Similarly, for confining groups $SO(N)$ $(Sp(N))$ with strictly real (resp., pseudo-real) representations the chiral symmetry $SU(n)$ of n Weyl fermions is likely to be broken down to the vector one $SO(n)$ (resp., pseudo-vector one $Sp(n)$) [12,17,18]. If so, the only candidates to be considered at all for composite model purposes are the non-Abelian gauge symmetries with complex (non-self-contragradient) representations.

It is well-known that, restricting oneself by the simple Lie groups, one encounters just three such possibilities: $SU(N), N \geq 3$; $SO(4k + 2), k \geq 2$ and exceptional group E_6

(see, e.g., [19]). The complex representations of the $SU(N)$ group can be anomaly free only if they contain necessarily higher rank tensors (in line with the fundamental ones, if desired). The $SO(4k+2)$ group, though being anomaly free, does not admit composite fermions built only of the constituent fermions in the fundamental (even dimensional spinor) representations.

On the other hand, E_6 group is free from both these drawbacks. First of all, E_6 is anomaly safe [19] in $d = 4$ dimensions so that there are no restrictions on its chiral fermion content. Besides, it possesses the odd (namely, the third) rank invariant tensor in the fundamental representation [20] and hence could lead to the required composite fermions.

Therefore one concludes that if one sticks to fermions in the fundamental representations of simple Lie groups, only chiral E_6 is permissible as S_{loc}. In this, the semi-simple Lie groups and/or nonfundamental representations, though not being excluded a priori, nevertheless seem quite unnatural for a truly underlying theory one is searching for.

Partial chiral symmetry breaking. A priori, for a strongly interacting gauge theory S_{loc} with chiral fermions there are two alternatives: either gauge symmetry is tumbled dynamically through its own strong interactions until all the constituent fermions are allowed to acquire dynamical masses, or the gauge symmetry remains exact and some of the chiral fermions have to remain massless. (In principle, some intermediate patterns could be adopted too.) It is the result of the dynamical competition between chiral symmetry breaking and confinement: which of these possibilities will win. Presently one does not know dynamical conditions under which either of them could be realized [18].

It is the second pattern (or at least some admixture of it) that is required for a composite picture of leptons and quarks to have any dynamical reason at all. Therefore, we take for granted that in the case under consideration underlying strongly coupled gauge symmetry E_6 is preserved, and proceed with studying the ensuing pattern of chiral symmetry breaking.

So, let for the chiral gauge theory S_{loc} the Lagrangian chiral symmetry G be dynamically broken to some vacuum residual symmetry $H : G \to H$. This is supposed to take place due to formation of vacuum bilinear condensate $< \chi_L \bar{\chi}_R >$ (plus $< \chi_R \bar{\chi}_L >$) from the constituent Weyl fermions $\chi_{L,R}$. We postulate that in this transition all those and only those constituents, which can get massive without breaking the confining gauge symmetry, do acquire dynamical masses. More than that, these masses are assumed to be equal. In other words, the hypothesis states that the residual chiral symmetry H is the maximal one consistent with the dynamical mass generation and preservation of the strongly coupled gauge symmetry. This agrees with the pattern of chiral symmetry breakdown adopted for vector-like, vector and pseudo-vector gauge theories, resp. $SU(N)$, $SO(N)$ and $Sp(N)$ [15]–[17]. Though this is just a hypothesis (a kind of the "survival" hypothesis) it is well-formulated and is more predictive than mere postulating some breaking pattern.

More explicitly, let in a general case of the chiral gauge E_6 symmetry there be l left-handed and r right-handed Weyl fermions χ_L and χ_R transforming as E_6 fundamental representation $\underline{N} = \underline{27}$. (Equivalently, in terms of left-handed fermions only, let there be l of $\underline{N}$'s and r of $\overline{N}$'s.) In general, $l \neq r$ are arbitrary. Asymptotic freedom of the gauge

E_6 requires only that $(l + r) < 22$. For definiteness let us assume that $l \geq r \geq 0$. The Lagrangian chiral symmetry G, left unbroken by the E_6 instantons, looks at different r as follows $(l \geq 2)$:

$$G = \begin{cases} SU(l)_L \times SU(r)_R \times U(1), \ r \geq 2; \\ SU(l)_L \times U(1), \ r = 1; \\ SU(l)_L, \ r = 0. \end{cases} \tag{1}$$

Under the hypothesis adopted, in the given case of chiral gauge E_6 the vacuum condensate in a suitably chosen basis can be brought to partly diagonal form as follows:

$$< \chi_L \bar{\chi}_R > = O(\Lambda_X^3) \begin{pmatrix} & & 0 & & \\ \cdots & \cdots & \cdots & \cdots & \cdots \\ 1 & & & 0 & \\ & \ddots & & & \\ & & 1 & & \\ & & & \ddots & \\ 0 & & & & 1 \end{pmatrix} \left. \begin{array}{c} \\ \\ \\ \end{array} \right\} l - r \quad \left. \begin{array}{c} \\ \\ \\ \\ \end{array} \right\} r \tag{2}$$

$$\underbrace{\qquad\qquad\qquad}_{r}$$

(Λ_X being the confinement mass scale of the exceptional gauge symmetry). This means that all r pieces of χ_R match some r pieces of χ_L leaving $n = l - r$ pieces of χ_L unmatched.

The condensate (Eq. 2) possesses the following residual symmetry H

$$H = \begin{cases} SU(n)_L \times SU(r) \times U(1), \ r \geq 2; \\ SU(n)_L \times U(1), \ r = 1; \\ SU(n)_L, \ r = 0, \end{cases} \tag{3}$$

where $n \equiv l - r$ is the net chirality index of the constituents ($0 \leq n \leq l$). (Here one should put $SU(n) \equiv I$ for $n = 0, 1$.) In the strictly chiral case ($r = 0$, $n = l$) the chiral symmetry is not broken at all ($G = H$), because according to the hypothesis the condensate just can not be formed without breaking the $S_{loc} = E_6$ gauge symmetry. In the other extreme vector-like case ($l = r$, $n = 0$) the condensate (Eq. 2) reduces in terms of Dirac fermions $\chi = (\chi_L, \chi_R)$ to the form $< \chi \bar{\chi} > \sim \mathrm{diag}(1, \ldots, 1)$, and the chiral symmetry is broken down to the vector-like one. In an intermediate case ($0 < n < l$) the chiral symmetry G is broken just partially (but to the maximum allowed extent). It is clear that all the constituents are in the vector-like representation $n \times \mathbf{1} \oplus \mathbf{r} \oplus \bar{\mathbf{r}}$ under the $SU(r)$ unbroken subgroup. The same can be shown to be true for the $U(1)$ residual subgroup.

Finally, we conclude that in the most general case of the chiral gauge E_6 the surviving chiral symmetry H is divided into two parts: strictly chiral and vector-like ones. Accordingly, there are two types of constituents relative to H: n massless Weyl fermions and r Dirac ones, the latter having equal dynamical masses. For the present purposes Dirac constituents are supposed to be intrinsically massless, though they could have some small explicit mass m ($m \ll \Lambda_X$).

5

Massless composite fermions. Rather general dynamical arguments require that chiral anomalies should match at the constituent and composite levels [3,14]. In this, anomalies for three unbroken currents have to match via massless composite fermions. Chiral anomaly matching condition is the unique known raison d' etre for the appearance of such massless states. Now we proceed to study this condition for residual subgroup H as given by Eq. 3.

It is well-known [20] that E_6 possesses totally symmetric invariant tensor d_{abc}; $a, b, c = 1, \ldots, 27$ in the fundamental representation, alongside with the Levi-Civita tensor $\epsilon_{abc\cdots}$, and so allows the formation of both three-particle and 27-particle fermion bound states. In what follows we restrict ourselves only with three particle composites.

In general, for chiral gauge E_6 there are three "strata" of composite fermions: pure chiral, mixed chiral and vector-like ones, built of two kinds of E_6 constituents, namely, strictly chiral and vector-like ones. Lorentz couplings of constituents have to be chosen in such a way as to allow for the formation of composite states of the required chiralities. A priori, left- and right-handed components of Dirac constituents enter these states in different Lorentz structures, in particular, those with derivatives. Hence, it is admitted that these constituents, though being potentially massive, in some chiral environment could not acquire their dynamical mass, so that corresponding composites are left massless. (In this respect massless chiral fermions somehow resemble composite Goldstone bosons, and for this reason one should think that potential models are not applicable to them.)

Now, let $\nu(\rho)$ be the chiral index of the state ρ (i.e. the number of the corresponding left-handed composite fermions minus that of the right-handed ones). $\nu(\rho)$ is some unknown integer which is supposed to be eventually determined by the underlying dynamics. The chiral anomaly matching condition can just somewhat restrict the allowable sets of these indices. Note that states composed exclusively of Dirac constituents should have zero indices due to the discrete LR-symmetry. Appropriate fermions fill in the vector-like composite stratum and have masses $O(\Lambda_X)$ (à la QCD hadrons).

The only triangle anomaly at the level of constituents is that $[SU(n)_L]^3$ for three $SU(n)_L$ currents. Solving the anomaly matching conditions one obtains in the most general case a three-parametric set of solutions. These solutions are, in general, not vector-like relative to the $SU(r) \times U(1)$ subgroup, though the constituents are. Nevertheless appropriate anomalies match and are equal to 0 in both cases.

Further reduction of the allowed set of indices could be achieved by imposing some additional physical restrictions. One of these is the matching condition of the mixed chiral-gravitational anomalies for one $U(1)$ and two gravitational currents [21]. This results in the requirement $\sum Y = 0$, where Y is the generator of the $U(1)$ subgroup of the unbroken chiral symmetry H. This gives one more relation for indices. The other possible restriction is decoupling condition [22] (in more refined form, persistent mass condition [23] or the constituent number independence [24]). But for chiral gauge theories the decoupling conditions are not obligatory [25].

In addition to massless composite fermions there also appear composite Goldstone bosons. They correspond to the broken symmetry G/H and are built of one chiral Weyl

and one vector-like Dirac constituents. Goldstone bosons saturate the anomaly matching for the broken currents from G/H.

Thus the chiral symmetry breaking pattern is just of the type required to embed the SM. But in order to built a particular realistic composite model based on this binding mechanism one has to specify a lot of "subtle" details, such as the quantum numbers of constituents, the intrinsic gauge symmetry, the explicit mass terms etc. Presently this cannot be done unambiguously.

Nevertheless the scheme do unambiguously produce the key message for the "low energy" effective theory of unified compositeness. Namely, it should be a nonlinear model G/H with G and H from the sets of Eq. 2 and Eq. 3, resp. This could be a starting point for studying the unified compositeness at the subthreshold energies. Additional topics of the scheme can be found in ref. [4].

3. Higgs doublet as composite Goldstone boson

Nonlinear Standard Model. To describe the "low energy" (i.e., below the compositeness scale) behaviour of the composite leptons, quarks and Higgs bosons, without detailed knowledge of the hyperstrong interactions responsible for their internal substructure, one has to refer to the framework of the effective field theory [9]. In essence, it requires just the assumption about the symmetry breaking pattern $G \to H$, as well as the light particle content. The simplest nonlinear model G/H to implement the idea of the Higgs doublet as composite Goldstone boson was first proposed on phenomenological grounds in ref. [2]. It was further refined from the unified compositeness point of view and systematically studied in ref. [5]. It may be called the minimal Nonlinear Standard Model (NSM). In what follows we present the basic features of the NSM. More details can be found in ref. [5].

It can be shown that the simplest nonlinear model G/H with the required properties is based on the symmetry breaking pattern $G = SU(3) \times U(1)$ and $H = I_{SM} = SU(2)_L \times U(1)_Y$. The extended symmetry G contains the broken isodoublet generators X_I, $X^{\dagger I}$, $I = 1, 2$, as well as the broken hypercharge Y' in addition to the unbroken SM generators of the weak isospin T_i, $i = 1, 2, 3$ and the weak hypercharge Y. With the broken generators of G/H there are associated Goldstone doublet ϕ_I and singlet ϕ'. This extra Goldstone boson is absorbed by the gauge boson of the additional dynamically broken local symmetry $U(1)_{Y'}$. The latter is the minimum one required to eventually convert the true Goldstone doublet ϕ via the radiative corrections into the SM Higgs doublet. In this prototype model, the QCD colour symmetry is supposed to be trivially present on both sides of the symmetry breaking chain.

Nonlinear realization. As the nonlinear model G/H, the NSM can be built via the canonical nonlinear realization of the symmetry G that becomes linear when restricted to H [10]. The Goldstone bosons parameterize the element of the left coset space $\xi \in G/H$

$$\xi = e^{i\phi'Y'/\mathcal{F}'} e^{i(\phi_I X^{\dagger I} + \text{h.c.})/\mathcal{F}}, \tag{4}$$

with $\mathcal{F}$, $\mathcal{F}' = \mathcal{O}(\mathcal{F})$ being the symmetry breaking mass scales. Here ξ and ϕ transform under $g \in G$ as

$$
\begin{aligned}
g \;:\; \xi \;&\to\; \tilde{\xi} = g\xi h^\dagger(g,\xi), \\
\phi \;&\to\; \tilde{\phi}(g,\xi),
\end{aligned}
\tag{5}
$$

where $h(g,\xi)$ and $\tilde{\phi}(g,\xi)$ are uniquely determined through the natural decomposition

$$
g\xi \equiv \tilde{\xi}h = e^{i\tilde{\phi}'Y'/\mathcal{F}'}\, e^{i(\tilde{\phi}_I X^{\dagger I}+\text{h.c.})/\mathcal{F}}h.
\tag{6}
$$

A matter field ψ transforms under $g \in G$ as

$$
g \;:\; \psi \;\to\; \rho(h(g,\xi))\psi,
$$

where ρ is a linear representation of H, and $h(g,\xi)$ is determined by the equation above.

Derivatives of the Goldstone and matter fields enter through the Maurer-Cartan 1-form $\Delta_\mu \equiv 1/i\,\xi^\dagger D_\mu \xi$, with D_μ being the derivative covariant w.r.t. the gauge symmetry I_{loc}. The 1-form Δ_μ contains the nonlinear covariant derivative $\mathcal{D}_\mu \phi$ of the Higgs-Goldstone doublet ϕ, as well as a part required to construct the nonlinear covariant derivative $\mathcal{D}_\mu \psi$ of the matter fields ψ. Namely, let us divide Δ_μ into two parts: $\Delta_{\|\mu}$ which is parallel to G/H and $\Delta_{\perp\mu}$ orthogonal to it, along the unbroken symmetry H:

$$
\begin{aligned}
\Delta_\mu \;&=\; (\Delta_{\|I\mu}X^{\dagger I} + \text{h.c.}) + \Delta^{0'}_{\|\mu}Y' \\
&+\; \Delta^i_{\perp\mu}T^i + \Delta^0_{\perp\mu}Y.
\end{aligned}
\tag{7}
$$

Then one has

$$
\begin{aligned}
(\mathcal{D}_\mu\phi)_I/\mathcal{F} \;&=\; \Delta_{\|I\mu}, \\
\mathcal{D}_\mu\psi \;&=\; \Big(\partial_\mu + i(\Delta^i_{\perp\mu}T^i + \Delta^0_{\perp\mu}Y)\Big)\psi.
\end{aligned}
\tag{8}
$$

All the terms in Eq. 7 transform nonlinearly under G as the irreducible representations of H and can be used to construct the effective Lagrangian of the NSM. It consists of the most general superficially H invariant expressions built of the ψ's (but not of ϕ's) and the nonlinear covariant derivatives $\mathcal{D}_\mu\phi$ and $\mathcal{D}_\mu\psi$, the latter ones transforming like ψ's under the nonlinearly realized extended symmetry G. Additional building blocks are given by the nonlinear generalization of the gauge field strengths $\Delta_{\mu\nu} \equiv 1/i\,\xi^\dagger[D_\mu, D_\nu]\xi$.

Let us decompose $\mathcal{L}_{eff}$ into the gauge, Higgs and fermion parts:

$$
\mathcal{L}_{eff} = \mathcal{L}_G + \mathcal{L}_H + \mathcal{L}_F.
\tag{9}
$$

Then the gauge part $\mathcal{L}_G$ of the Lagrangian is built of the irreducible under H components of $\Delta_{\mu\nu}$, the latter ones being defined as in Eq. 7. The Higgs part

$$
\mathcal{L}_H = (\mathcal{D}_\mu\phi)^\dagger(\mathcal{D}_\mu\phi) + \mathcal{O}(1/\mathcal{F}^4)
\tag{10}
$$

8

is uniquely determined by the symmetry breaking pattern. And finally, the fermion part for the left-chiral fermion fields ψ_ρ, belonging to the irreducible representation ρ of the unbroken subgroup H, takes the form

$$
\begin{aligned}
\mathcal{L}_F &= \sum_\rho (m_\rho \psi_{\bar\rho} \psi_\rho + \text{h.c.}) + \sum_\rho \bar\psi_\rho \sigma_\mu i/2 \overset{\leftrightarrow}{\mathcal{D}}_\mu \psi_\rho \\
&+ \sum_\rho \eta_\rho \bar\psi'_\rho \sigma_\mu \psi_\rho \Delta^{0\prime}_{\|\mu} + \text{h.c.} \\
&+ \frac{1}{\mathcal{F}} \sum_{\rho_1 \rho_2} \chi_{\rho_1 \rho_2} \bar\psi^I_{\rho_2} \sigma_\mu \psi_{\rho_1} \mathcal{D}_\mu \phi_I + \text{h.c.} \\
&+ \frac{1}{\mathcal{F}} \sum_{\rho_1 \rho_2} \bar\chi_{\rho_1 \rho_2} \bar\psi^I_{\rho_2} \sigma_\mu \psi_{\rho_1} \overline{(\mathcal{D}_\mu \phi)}_I + \text{h.c.} + \mathcal{O}(\frac{1}{\mathcal{F}}).
\end{aligned} \tag{11}
$$

We have omitted terms irrelevant for the further discussion. Here m, η, χ and $\bar\chi$ are arbitrary parameters. The first expression in $\mathcal{L}_F$ describes the explicit mass terms of the vector-like heavy composite fermions ($m_\rho = \mathcal{O}(\mathcal{F})$). It can be shown that all the terms $\mathcal{O}(1/\mathcal{F})$ mix with necessity the light chiral and heavy vector-like fermions. In the limit $\mathcal{F} \to \infty$ the finite part of $\mathcal{L}_{eff}$ reproduces exactly the SM Lagrangian (except for the Higgs potential and Yukawa interactions). In this limit the heavy vector-like fermions decouple from the SM light sector.

Higgs and Yukawa interactions. In reality, symmetry G is not exact but explicitly violated, e.g., by the extended electroweak interactions, since only part of G, namely $I_{loc} = SU(2)_L \times U(1)_Y \times U(1)_{Y'}$, is supposed to be gauge. Gauge radiative corrections may lead to a misalignment of the dynamically unbroken subgroup H relative to the gauge I_{loc}^{SM}. This results in the spontaneous SM symmetry breaking and the appearance of the Higgs and Yukawa effective interactions. This effect may be properly accounted for by adding the symmetry violating effective Lagrangian

$$
\Delta\mathcal{L}_H = \mathcal{F}^4 \Big(\bar g^2 tr(\xi^\dagger T_i \xi T_i) + \bar g_1^2 tr(\xi^\dagger Y \xi Y) - \bar g_1'^2 tr(\xi^\dagger Y' \xi Y') \Big). \tag{12}
$$

Here the effective couplings $\bar g^2$, $\bar g_1^2$ and $\bar g_1'^2$ are equal to the product of the corresponding gauge constants squared and some spectral integrals. Note the difference in the sign between the contributions of the dynamically broken and unbroken gauge interactions [11, 12]. Decomposition of $V_H = -\Delta\mathcal{L}_H$ in the region of weak fields ($|\phi|/\mathcal{F} \ll 1$) gives the Higgs potential up to $\mathcal{O}(1/\mathcal{F}^2)$. Note that the Higgs boson is naturally expected to be light in the scheme.

As for the Yukawa interactions, $\mathcal{L}_{eff}$ includes three ingredients required for their appearance: the chirality changing mass terms of the heavy vector-like fermions ($m = \mathcal{O}(\mathcal{F})$), the Goldstone interactions of these fermions ($\sim \mathcal{D}_\mu \phi/\mathcal{F}$), and, finally, the weak gauge mixing of the light chiral and heavy vector-like fermions. So, the loop corrections may lead to the appearance of the symmetry violating effective Lagrangian like Eq. 12. Its decomposition can be shown to result in the nonderivative Yukawa couplings of order $\mathcal{O}(g^2/(4\pi)^2)$. More details can be found in ref. [5].

4. Vector boson dominance of gauge interactions

Hidden local symmetry. Being a nonlinear model G/H, the NSM is equivalent to the model with linearly realized symmetry $G \times \hat{H}_{loc}$ [26]. Here $\hat{H}_{loc} \simeq H$ is the hidden local symmetry of the original NSM with the appropriate auxiliary gauge bosons. In the context of the minimal NSM the phenomenon of the hidden local symmetry has been first studied in ref. [6]. The essence of the latter one is as follows.

In the linear model, the field variable is the element of the whole group G that can be parameterized as

$$\hat{\xi} = \xi h, \ \ h \in H. \tag{13}$$

The following transformation law under $g \times \hat{h}(x) \in G \times \hat{H}_{loc}$ takes place:

$$g \times \hat{h}(x): \ \hat{\xi} \to g\hat{\xi}\hat{h}^\dagger(x). \tag{14}$$

The linear model describes dynamical/spontaneous symmetry breaking $G \times \hat{H}_{loc} \to H$, with the total local symmetry being broken as $I_{loc} \times \hat{H}_{loc} \to I_{loc}^{SM} = SU(2)_L \times U(1)_Y$.

To construct the Lagrangian of the linear model one has to introduce the modified 1-form $\hat{\Delta}_\mu = 1/i\,\hat{\xi}^\dagger \hat{D}_\mu \hat{\xi}$, with $\hat{D}_\mu$ being now the derivative covariant both under the intrinsic gauge symmetry I_{loc} and the hidden local symmetry $\hat{H}_{loc}$. Let us again divide $\hat{\Delta}_\mu$ into two parts: $\hat{\Delta}_{\|\mu}$ and $\hat{\Delta}_{\perp\mu}$. Under $G \times \hat{H}_{loc}$ the longitudinal part $\hat{\Delta}_{\|\mu}$ transforms homogeneously as in the original nonlinear model, and so does now the transversal part $\hat{\Delta}_{\perp\mu}$. It is precisely the auxiliary vector fields $\hat{W}_\mu^i$ and $\hat{S}_\mu$, corresponding to $\hat{H}_{loc}$ which make the transformation of $\hat{\Delta}_\perp$ homogeneous. In the unitary under $\hat{H}_{loc}$ gauge, i.e. at $h \equiv 1$ in Eq. 13, the modified 1-form looks like

$$\begin{aligned} \hat{\Delta}_{\|\mu} &= \Delta_{\|\mu}, \\ \hat{\Delta}_{\perp\mu}^i &= \Delta_{\perp\mu}^i - \hat{g}\hat{W}_\mu^i, \\ \hat{\Delta}_{\perp\mu}^0 &= \Delta_{\perp\mu}^0 - \hat{g}_1\hat{S}_\mu, \end{aligned} \tag{15}$$

where Δ_μ is the 1-form present in the original minimal NSM, $\hat{g}$ and $\hat{g}_1$ being some new strong coupling constants (supposedly, $\hat{g}^2/4\pi = \mathcal{O}(1)$).

In the effective Lagrangian of the linear model, the new terms appear. They are related to the orthogonal part of the modified 1-form. Here are some of the appropriate terms in the gauge sector:

$$\mathcal{L}_G = \frac{\lambda\mathcal{F}^2}{2}(\hat{\Delta}_{\perp\mu}^i)^2 + \frac{\lambda_1\mathcal{F}^2}{2}(\hat{\Delta}_{\perp\mu}^0)^2 + \cdots, \tag{16}$$

and for the chiral fermions they are

$$\begin{aligned} \mathcal{L}_F &= \bar{\psi}\sigma_\mu i(\partial_\mu + i\hat{g}\hat{W}_\mu^i T^i + i\hat{g}_1\hat{S}_\mu)\psi \\ &+ \kappa\bar{\psi}\sigma_\mu T^i\psi\hat{\Delta}_{\perp\mu}^i + \kappa_1\bar{\psi}\sigma_\mu Y\psi\hat{\Delta}_{\perp\mu}^0 + \cdots. \end{aligned} \tag{17}$$

Here λ's and κ's are free parameters. It is to be noted that the matter fields ψ transform now only under $\hat{H}_{loc}$. The modified covariant derivative for them contains only the

composite $\hat{W}_\mu$ and $\hat{S}_\mu$, but not the elementary W_μ and S_μ, the latter ones entering only through the nonminimal interactions.

Introducing the vector fields in such a way without kinetic terms is just a formal procedure. But we believe that the required kinetic terms are developed by the quantum effects, and the new composite vector bosons become physical. This takes place, e.g., in 2- and 3-dimensional nonlinear σ-models [27], as well as in the hadron physics as an accomplished fact.

Vector boson dominance. From the Lagrangian of the linear model, one can read off the Lagrangian terms of the vector boson-current interactions:

$$
\begin{aligned}
\mathcal{L}_{int} &= -gW_\mu^i\Big((1-\lambda)J_\mu^i(\phi) + \kappa J_\mu^i(\psi)\Big) \\
&\quad -\hat{g}\hat{W}_\mu^i\Big(\lambda J_\mu^i(\phi) + (1-\kappa)J_\mu^i(\psi)\Big).
\end{aligned}
\tag{18}
$$

Here $J_\mu^i(\psi) = \bar{\psi}\gamma_\mu T^i\psi$ and $J_\mu^i(\phi) = \phi^\dagger i\tau^i/2 \overset{\leftrightarrow}{D}_\mu \phi$ are the usual SM isotriplet currents, with D_μ being the SM covariant derivative. To these isospin terms, one has to add the similar hypercharge isosinglet terms. Impose now the natural requirement that all the composite particles ϕ and ψ interact directly only with the composite vector bosons $\hat{W}$ and $\hat{S}$, but not with the elementary ones W and S. In other words, this is the well-known hypothesis of the vector boson dominance (VBD). This requirement allows one to fix the free parameters: $\lambda = 1$, $\kappa = 0$ and similarly for the isosinglet parameters.

The terms $(\hat{\Delta}_\perp^i)^2$ and $(\hat{\Delta}_\perp^0)^2$ describe the mass mixing of the elementary and composite gauge bosons, namely, W with $\hat{W}$ and S with $\hat{S}$. Diagonalizing these terms one gets two sets of physical vector bosons: the massless isotriplet and isosinglet physical bosons $\bar{W}^i$ and $\bar{S}$, as well as the massive ones $\bar{\hat{W}}^i$ and $\bar{\hat{S}}$ with masses of order $\mathcal{F}$. Due to the heavy physical vector boson exchange, the new low energy effective current-current interactions appear in addition to that of the SM:

$$
\begin{aligned}
\mathcal{L}_{int}^{(VBD)} &= -\frac{1}{2\mathcal{F}^2}\Big(J_\mu^i(\psi)J_\mu^i(\psi) + \eta_1 J_\mu^0(\psi)J_\mu^0(\psi)\Big) \\
&\quad -\frac{1}{\mathcal{F}^2}\Big(J_\mu^i(\psi)J_\mu^i(\phi) + \eta_1 J_\mu^0(\psi)J_\mu^0(\phi)\Big).
\end{aligned}
\tag{19}
$$

Here η_1 is a free parameter, related to the original minimal NSM. Note that the VBD does not affect the low energy Higgs boson self-interactions, the latter ones being determined by the original minimal NSM alone:

$$
\mathcal{L}_{int}(\phi) = -\frac{1}{\mathcal{F}^2}\Big(\frac{1}{3}J_\mu^i(\phi)J_\mu^i(\phi) + J_\mu^0(\phi)J_\mu^0(\phi)\Big).
\tag{20}
$$

All these expressions are valid only at energies $\sqrt{s} \ll \mathcal{F}$.

To resume, the unified compositeness plus the VBD prescribe the two-parameter set of the universal residual fermion-fermion, fermion-boson and boson-boson interactions, with their space-time and internal structure being fixed including sign. The unified compositeness scale $\mathcal{F}$ is expected to be in the deca-TeV region. Hence, the TeV energies are required to probe these new contact interactions.

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5. Universal dominant residual interactions

VBD of electroweak interactions. We have investigated the potential to test the hypothesis of the VBD of electroweak interaction at the future 2 TeV e^+e^- linear collider via $e^+e^- \to \bar{f}f$ [7] and $e^+e^- \to ZH$, W^+W^- [8]. We chose for studying a set of integral characteristics: the relative deviation Δ in the total cross-sections from the SM values, the forward-backward charge asymmetry A_{FB}, the left-right polarization asymmetry A_{LR} and the mixed asymmetry A_{LR}^{FB}.

We have calculated these observables for the processes $e^+e^- \to \mu^+\mu^-$ ($\tau^+\tau^-$), $\bar{b}b$, $\bar{c}c$, *jet jet* and for the Bhabha scattering $e^+e^- \to e^+e^-$ as functions of the parameter η_1 for the various values of $\mathcal{F}$. The general results of these calculations are as follows. For all the processes (except Bhabha scattering) all the asymmetries have the similar behaviour. First of all, there exists a particular value of $\eta_1 = \tan^2\theta_W \simeq 0.3$ when all the asymmetries coincide with those of the SM. The only way to unravel the contact interactions in this particular case is to study directly the total cross-sections. Another particular value of $\eta_1 = g_1^2\mathcal{F}^2/s$ provides the best case for studying the contact interactions, when all the asymmetries in all the processes saturate their maximal values.

To evaluate the statistical significance of the observed deviations we have considered the total cross-sections. Fig. 1 presents the reach for the scale $\mathcal{F}$ at 2σ level (95% C.L.) via the total cross-sections in the various $\bar{f}f$ channels.
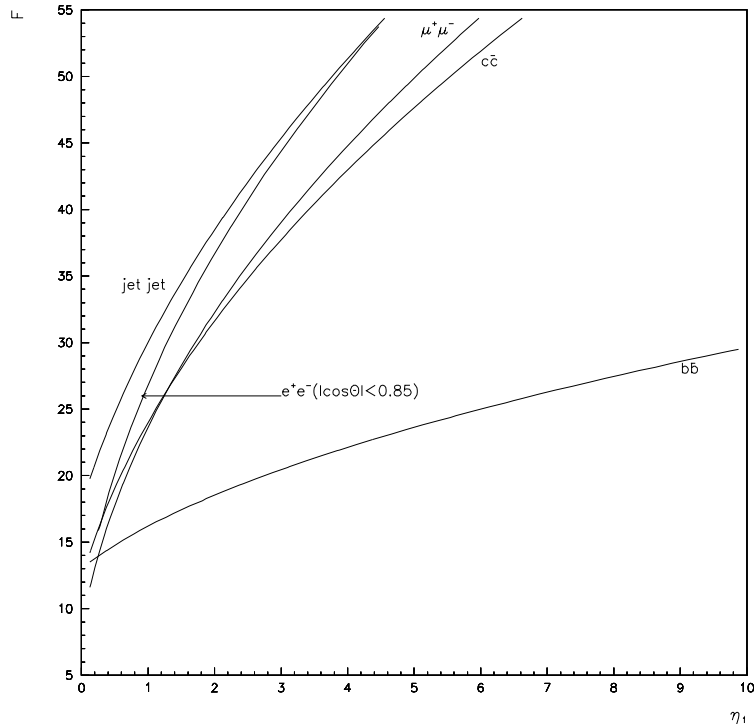

Fig. 1. The reach at 95% C.L. for the compositeness scale $\mathcal{F}$, vs. the parameter η_1, via studying the total cross-sections of the processes $e^+e^- \to \bar{f}f$.

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To this end we took into account only the statistical errors and accepted the integrated luminosity $\int \mathcal{L}\, dt$ moderately to be 20 fb^{-1}. In the case of the Bhabha scattering $e^+e^- \to e^+e^-$ an optimal value of the cutoff, equal to 0.85, was chosen. Here the sensitivity is maximal due to the maximal suppression of the t-channel peak at the statistics still high enough. It is seen that in the processes $e^+e^- \to \bar{f}f$ the VBD can be tested for the unified substructure scale $\mathcal{F}$ up to $\mathcal{O}(50\ \text{TeV})$.

For the processes $e^+e^- \to ZH$ and W^+W^-, it proved to be of importance to consider the polarized cross-sections $\sigma(P_e)$, with P_e denoting the polarization of electron beam (the positron beam was taken to be unpolarized). So, we have studied the relative deviation $\Delta(P_e)$ in the polarized cross-section from that of the SM. In the cases of both ZH and WW pair production one has $|\Delta(-1)| \ll |\Delta(+1)|$. Hence one is lead to conclude that it is preferable to operate with the maximum right-handedly polarized electrons to observe as large deviations in the total cross-sections from the SM values as possible. The advantage of the right-handed polarization can be seen, e.g., from the picture that presents the scale $\mathcal{F}$ versus the parameter η_1, attainable at 95% C.L. (Fig. 2).
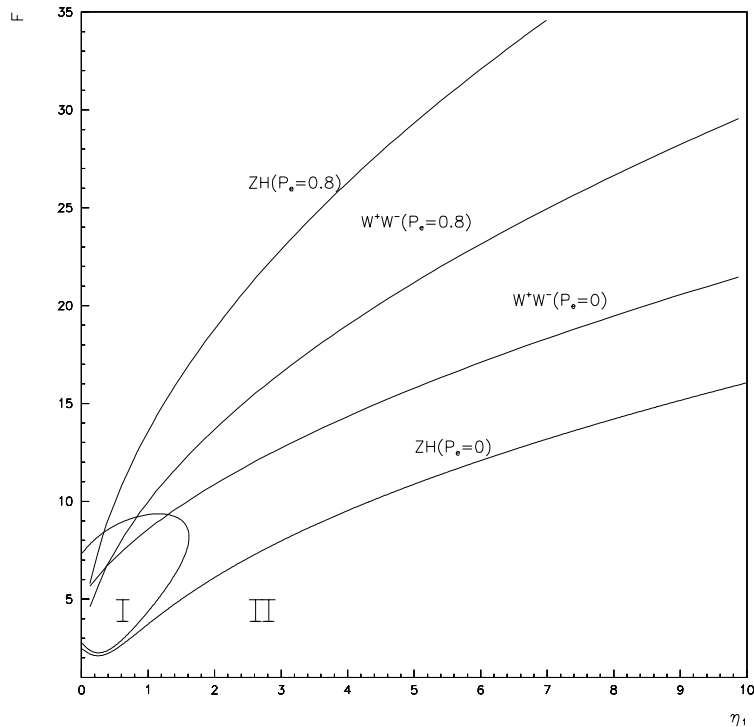

Fig. 2.   The same as in Fig. 1 for the processes $e^+e^- \to ZH$, $^+W^-$ with the various electron polarizations P_e ($m_H = 200$ GeV).

Thus, using the right-handed polarized electron beam the VBD can be tested up to the scale $\mathcal{F}$ of the order of 25 TeV in the e^+e^- annihilation into boson pairs. Here the calculations for the W^+W^- pair production have been made under the instrumental cutoff $|\cos\theta| \leq 0.8$. In addition, an optimal cutoff in the forward direction, whose sense is similar to that in the forward Bhabha scattering, has been found to be $\cos\theta \leq 0.3$.

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Anomalous triple gauge interactions. In addition to the VBD interactions, a lot of other "low energy" residual interactions is allowed in the scheme of the unified compositeness. In particular, the exotic triple gauge interactions (TGI) [28] are conceivable too, and can contribute to the W^+W^- pair production. The question arises as to what extent the two types of new interactions could imitate each other.

The anomalous TGI should originate from a kind of the SM extension. Here, the SM symmetry $SU(2)_L \times U(1)_Y$ could be realized either linearly or nonlinearly. In the case of the nonlinear realization (being still linear on the $U(1)_{em}$ subgroup), the nonlinearity scale Λ is just the SM v.e.v. v. Thus, this kind of extension has nothing to do with the unified compositeness we consider. On the other hand, for the linear SM symmetry realization the scale Λ is not directly related with v and could be as high as desired. Thus, we chose it to be the unified compositeness scale $\mathcal{F} = \mathcal{O}(10 \text{ TeV})$.

All the conceivable linearly realized residual interactions are described by the $SU(2)_L \times U(1)_Y$ invariant operators built of the SM fields [29,30]. All the operators which are relevant to the anomalous TGI vertices are naturally expected to be $\mathcal{O}(g)$ or less in the gauge couplings, but one exception, namely, $\mathcal{O}_{WS}$. The latter stems from the nonlinear generalization of the field strengths in the NSM. The similar gauge kinetic terms of the isotriplet W and isosinglet S bosons have no gauge couplings. So, the same must naturally happen for $\mathcal{O}_{WS}$, for its origin is of the same nature.

Thus, we have retained the $\mathcal{O}_{WS}$ operator alone and have chosen the proper effective Lagrangian to be

$$\mathcal{L}_{eff} = \frac{C}{2}\frac{1}{\mathcal{F}^2}\mathcal{O}_{WS} \equiv \frac{C}{2}\frac{1}{\mathcal{F}^2}\phi^\dagger\frac{\tau_i}{2}\phi W_{\mu\nu}^i S_{\mu\nu}, \tag{21}$$

where $C = \mathcal{O}(1)$. With account for all the contributions from this operator we have found that the deviations from the SM predictions even in the most enhanced TGI case are much smaller than those in the VBD case. So, the VBD is in fact dominant.

Conclusion

The scenario of unified compositeness of leptons, quarks and Higgs bosons, with the unification of the Higgs and Yukawa interactions as residual ones, is the viable alternative to presently popular scenarios of New Physics with the elementary point-like fields and fundamental interactions. This scenario allows one to have a fresh look at the old problems and to put forward the new ones. The naturally preferred Deca-TeV compositeness scale makes the scenario amenable to experimental study at the future TeV energy colliders. If realized in Nature, this scenario would open completely new prospects for the whole high energy physics development.

References

[1] M.E.Peskin, in *Proc. of the Int. Symp. on Lepton and Photon Interactions at High Energies*, Bonn, 1981, p. 880.

[2] D.B.Kaplan and H.Georgi, *Phys. Lett.* **136B**, 183 (1984); H.Georgi and D.B.Kaplan, *Phys. Lett.* **145B**, 216 (1984); M.Dugan, H.Georgi and D.B.Kaplan, *Nucl. Phys.* **B254**, 299 (1985).

[3] G.'t Hooft, in *Recent Developments in Gauge Theories. Proc. of the NATO Advanced Study Inst.*, Cargèse, 1979, eds. G.'t Hooft et al. (Plenum Press, New York, 1980) p. 135.

[4] Yu.F.Pirogov, *Int. J. Mod. Phys.* **A9**, (1994).

[5] Yu.F.Pirogov, *Int. J. Mod. Phys.* **A7**, 6473 (1992); in *Proc. of the Int. Conf. "Quarks '92"*, Zvenigorod, 1992, eds. D.Yu.Grigoriev *et al.* (World Scientific, Singapore, 1993) p. 375.

[6] Yu.F.Pirogov, *Mod. Phys. Lett.* **A8**, 3129 (1993); in *Proc. of the Int. School "Particles and Cosmology"*, Baksan Valley, 1993, eds. E.N.Alexeev *et al.* (World Scientific, Singapore, 1994) p. 151.

[7] V.V.Kabachenko and Yu.F.Pirogov, *Int. J. Mod. Phys.* **A10**, 3187 (1995).

[8] V.V.Kabachenko and Yu.F.Pirogov, preprint IHEP 95-92, 1995; to be publ. in *Int. J. Mod. Phys.* **A11** (1996).

[9] S.Weinberg, *Physica* **96A**, 327 (1979).

[10] S.Coleman, J.Wess and B.Zumino, *Phys. Rev.* **177**, 2239 (1969); C.G.Callan, S.Coleman, J.Wess and B.Zumino, *ibid*, 2247.

[11] R.Dashen, *Phys. Rev* **D3** 1879 (1971); J.P.Preskill, *Nucl. Phys.* **B177** 21 (1981); T.Banks, *Nucl. Phys.* **B243** 125 (1984).

[12] M.E.Peskin, *Nucl. Phys.* **B175** 197 (1980).

[13] D.A.Cosower, *Nucl. Phys.* **B294**, 244 (1987); V.Baluni, *Phys. Lett.* **222B**, 103 (1989).

[14] Y.Frishman *et al.*, *Nucl. Phys.* **B177** 157 (1981); S.Coleman and B.Grossman, *Nucl. Phys.* **B203** 205 (1982).

[15] S.Coleman and E.Witten, *Phys. Rev. Lett.* **45**, 100 (1980).

[16] D.Weingarten, *Phys. Rev. Lett.* **51** 1830 (1983); S.Nussinov, *ibid.* 2083; E.Witten, *ibid.* 2351.

[17] S.Dimopoulos, *Nucl. Phys.* **B168** 69 (1980).

[18] M.E.Peskin, in *Recent Advances in Field Theory and Statistical Mechanics. Les Houches Summer School of Theoretical Physics*, 1982, eds. J.Zuber and R.Stora (North-Holland, Amsterdam, 1984).

[19] S.Okubo, *Phys. Rev.* **D16** 3528, 3535 (1977).

[20] P.Cvitanović, *Phys. Rev.* **D14** 1536 (1976).

[21] R.Delbourgo and A.Salam, *Phys. Lett.* **40B** 381 (1972); T.Eguchi and P.Freund, *Phys. Rev. Lett.* **37** 1251 (1976); L.Alvarez-Gaume and E.Witten, *Nucl. Phys.* **B234** 269 (1983).

[22] K.Symanzik, *Comm. Math. Phys.* **34** 7 (1973); T.Appelquist and J.Carazzone, *Phys. Rev.* **D11** 2856 (1975).

[23] J.P.Preskill, S.Weinberg, *Phys. Rev.* **D24** 1059 (1981).

[24] R.Kaul and R.Rajaraman, *Phys. Lett.* **110B** 385 (1982).

[25] M.S.Chanowitz and M.K.Gaillard, *Nucl. Phys.* **B261**, 379 (1985).

[26] M.Bando, T.Kugo, S.Uehara, K.Yamawaki and T.Yanagida, *Phys. Rev. Lett.* **54**, 1215 (1985); M.Bando, T.Kugo and K.Yamawaki, *Nucl. Phys.* **B259**, 493 (1985); *Prog. Theor. Phys.* **73**, 1541 (1985); *Phys. Rep.* **164**, 217 (1988).

[27] A.D'Adda, P. di Vecchia and M.Lüscher, *Nucl. Phys.* **B146**, 63 (1978); *Nucl. Phys.* **B152**, 125 (1979); A.M.Polyakov, Gauge Fields and Strings, *Contemporary Concepts in Physics*, v. 3 (Harwood Academic Publishers, 1987) p. 139.

[28] K.Hagiwara, R.D.Peccei, D.Zeppenfeld and K.Hikasa, *Nucl. Phys.* **B282**, 253 (1987); M.A.Samuel *et al.*, *Phys. Rev. Lett.* **67**, 9 (1991).

[29] C.N.Leung, S.L.Love and S.Z.Rao, *Z. Phys.* **C** — *Part. and Fields* **31**, 433 (1986); W.Buchmuller and D.Wyler, *Nucl. Phys.* **B236**, 621 (1986).

[30] A.de Rujula, M.B.Gavela, P.Hernandes and E.Masso, *Nucl. Phys.* **B384**, 3 (1992).

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