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STANDARD MODEL WITHOUT HIGGS PARTICLES

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1 Introduction

Many problems of the conventional standard model (SM) arise from the presence of Higgs sector of scalar fields. Perhaps one of the most evident problems is the absence of any experimental manifestations of Higgs particles. The ordinary explanation of this fact by a large Higgs boson mass M_H may be unsatisfactory. At fairly large M_H, approximately at $M_H \geq 1$ TeV [1], [2], the SM becomes a strongly interacting theory. In this case the usual perturbation theory (PT), which is the only reliable method of calculation in quantum field theory (QFT), cannot be applied to derivation of the SM predictions. In the near future at the SSC and LHC the Higgs mass range $M_H \leq 800$ GeV will be exhaustively explored [3], [4] and the above-mentioned upper bound of the perturbative regime of the SM may be exceeded.

A situation when Higgses are not discovered is now under serious theoretical consideration. The composite models and schemes with dynamical mechanisms of gauge symmetry breaking are investigated extensively as an alternative of the models with the fundamental scalars. General properties of spontaneously broken gauge symmetry (SBGS) are also explored independently of concrete symmetry breaking mechanisms [2], [5]. To our mind, it is interesting to consider other approaches to SBGS without observable Higgs particles in the framework of QFT.

We propose an approach based on the idea of the so-called "virton" field [6] which can be constructed within the nonlocal QFT. In the local one it doesn't exist. The main peculiarity of the virton field is the following. After quantization it describes not ordinary particles, but unobservable quasi-particles, which appear only as virtual states. If we regard the virton fields as the Higgs fields, we obtain an appropriate model of the Higgs sector of SM which generates the vector boson and fermion masses without producing observable Higgs particles.

The nonlocal quantum field theory (NLQFT) is a self-consistent scheme satisfying all principles of the conventional QFT (unitarity, causality, relativistic invariance, etc.) and providing the basis for correct description of the nonlocality effects. At the same time the nature of nonlocality itself may be unknown. This point of view is accepted in the series of work [7] where a version of SM is developed in which all interactions are nonlocal. There are some problems with the gauge invariance because of nonlocality in gauge field interactions. Nevertheless, this approach gives a good possibility of building completely ultraviolet finite theory of the fundamental interactions. Another way to the finite QFT was proposed in the work [8] on the basis of infinite component fields which also results in the introduction of the special form of nonlocality.

In our approach the nonlocality is introduced only in the Higgs self-interaction term. The main goal of this modification of the SM is to exclude the scalar particles from the observable spectrum. The theory in this case is certainly not finite, though its ultraviolet property is improved and divergences of many diagrams are reduced. Our method of quantization of the nonlocal fields also differs from the one applied to this problem in the above mentioned approach [7]. We use intermediate regularization of the nonlocal field theory by introduction the infinite set of quantized local auxiliary fields defined on the Hilbert space with the negative-norm states.

2 Electroweak Symmetry Breaking and Nonlocal Self Interaction of Higgs Fields

We introduce nonlocality into the Higgs self-interaction writing down the Lagrangian of the scalar electroweak doublet fields in the form

$$-\mathcal{L} = \phi^\dagger(x)(\partial^2 + m^2)\phi(x) + \lambda\left(\Phi(x)^\dagger * \Phi(x)\right)^2 \tag{1}$$

where $m^2 < 0$, and the nonlocal field $\Phi(x)$ is obtained from the local one $\phi(x)$ by "smearing" over the nonlocality domain with the characteristic scale ℓ_0. We don't specify the nature of this nonlocality and, introducing the phenomenological formfactor $\mathcal{K}$, define the nonlocal field

$$\Phi(x) = \int dy\, \mathcal{K}(x-y)\, \phi(y) = \mathcal{K}\left(\ell_0^2\, \partial^2\right)\, \phi(x) \tag{2}$$

The nonlocal operator $\mathcal{K}(\ell_0^2\, \partial^2)$ can be presented in the form

$$\mathcal{K}\left(\ell_0^2\, \partial^2\right) = \sum_{n=0}^{\infty} \frac{c_n}{(2n)!}\left(\ell_0^2\, \partial^2\right)^2 \tag{3}$$

Then the generalized function $\mathcal{K}(x-y) = \mathcal{K}(\ell_0^2\, \partial^2)\, \delta(x-y)$ belongs to one of the spaces of nonlocal generalized functions which was introduced and explored in the works of Efimov [9].

Considering the theory based on the Lagrangian (1) we follow the method of quantization of the nonlocal fields developed in these works.

Let us rewrite the Lagrangian (1) in terms of the nonlocal fields $\Phi(x)$

$$-\mathcal{L} = \Phi^\dagger(x)\mathcal{K}^{-2}\left(\ell_0^2\, \partial^2\right)\left(\partial^2 + m^2\right)\Phi(x) + \lambda\left(\Phi(x)^\dagger * \Phi(x)\right)^2 \tag{4}$$

We are looking for such conditions which, being applied to this Lagrangian, guarantee the virton realization for the scalar field remaining after SBGS. In this case observable scalar particles will not appear because their propagator $G(p^2)$ is an entire function. Introduce a new operator

$$\mathcal{E}\left(\partial^2\right) = \mathcal{K}^{-2}\left(\ell_0^2\, \partial^2\right)\left(\partial^2 + m^2\right) + \omega \tag{5}$$

The constant ω will be fixed from the condition $G(p^2) = \mathcal{E}^{-1}(-p^2)$. Then we also require that the function $\mathcal{K}$ is an entire analytic function without any zeros. This means that the Higgs propagator $G(p^2)$ has no poles after SBGS. The $\mathcal{E}$-function must satisfy some conditions following from the general principles of QFT. These are the Efimov conditions [9]

- $\mathcal{E}(z)$ -is an entire analytic function of the order $\frac{1}{2} \leq \rho \leq 1$,
- $[\mathcal{E}(z)]^* = \mathcal{E}(z^*)$
- $\mathcal{E}(z) > 0$ for real z
- In Euclidean momentum space $\mathcal{E}^{-1}(-p^2)$ has to decrease steeply enough for

$$\int_0^\infty dp_E^2\ \mathcal{E}^{-1}(p_E^2) < \infty \tag{6}$$

The last condition results in decreasing Euclidean Green functions of NLQFT. The most general form of the $\mathcal{E}$-function satisfying all these conditions is

$$\mathcal{E}(z) = \mu^2 \exp\left(\mathcal{W}(z)\right) \tag{7}$$

where μ is a parameter providing $\mathcal{E}(z)$ with the correct dimension ; $\mathcal{W}$ is a real entire function increasing with $z^2 \to \infty$.

The interaction with the gauge fields is introduced by usual minimal substitution

$$\partial_\mu \longrightarrow D_\mu = \partial_\mu - ig\frac{\tau^a}{2}\mathsf{A}_\mu^a - ig'\frac{Y}{2}\mathsf{B}_\mu \tag{8}$$

where A_μ and B_μ are the SU_{2L} and U_{1Y} gauge fields, respectively; Y is a weak hypercharge operator.

Taking these into account, we rewrite the Lagrangian (4) in the form

$$-\mathcal{L} = \Phi^\dagger(x)\left(\mathcal{E}\left(D^2\right) - \omega\right)\Phi(x) + \lambda\left(\Phi(x)^\dagger * \Phi(x)\right)^2 \tag{9}$$

Gauge invariance of this Lagrangian is the direct consequence of the fact that the $\mathcal{E}(z)$-function is an entire one. Under gauge transformations

$$\Phi \xrightarrow{g} \Phi^g = g\Phi \tag{10}$$

$$D_\mu \xrightarrow{g} D_\mu^g = g D_\mu g{-1} \tag{11}$$

the operator $\mathcal{E}(D^2)$, being the sum of positive degrees of the covariant derivative D, transforms as

$$\mathcal{E} \xrightarrow{g} \mathcal{E}^g = \mathcal{E}(gD^2g^{-1}) = g\mathcal{E}(D^2)g^{-1} \tag{12}$$

Noteworthy is that according to the Pikar theorem on a - points of an entire function [10], the $\mathcal{E}(z)$ function takes the ω-value an infinite number of times. This statement can be easily understood with the simplest example of function $\exp z$. Consequently, the propagator $\left(\mathcal{E}(-p^2) - \omega\right)^{-1}$ of the Φ – field has an infinite number of poles, some of them at negative or imaginary values of p^2. In quantum theory they correspond to particles with unphysical complex masses. As is well known, the presence of these states, analogous to "tachyon" states in the conventional SM, are the signal of SBGS.

Let us consider the quantization problem for the theory based on the Lagrangian (9). The standard canonical procedure cannot be directly applied in this case due to the presence of higher time derivatives in the kinetic term. Therefore, following [9], we consider the properly regularized theory

$$-\mathcal{L}^\delta = \Phi^{\delta\dagger}(x)\mathcal{E}^\delta\left(D^2\right)\Phi^\delta(x) - \omega^\delta\Phi^{\delta\dagger}(x) * \Phi^\delta(x) + \tag{13}$$

$$+ \lambda\left(\Phi^{\delta\dagger}(x) * \Phi^\delta(x)\right)^2 \tag{14}$$

where δ is a regularization parameter. The regularization is chosen in such a way that

$$\lim_{\delta\to 0}\mathcal{E}^{\delta}\left(D^2\right) = \mathcal{E}\left(D^2\right) \tag{15}$$

and the regularized function

$$\mathcal{E}^{\delta}\left(p^2\right) \sim \prod_{j=1}^{\infty}\left(p^2 - m_j^2(\delta)\right) \tag{16}$$

has an infinite number of zeros in a sequence of points such that $m_j^2(\delta) > 0$ and $m_j^2(\delta) \to \infty$ when $\delta \to 0$.

Then the theory can be quantized on the Hilbert space with the negative-norm states [11].

For the $\delta > 0$ the free Hamiltonian H_0^δ, the S^δ–matrix, the Green functions G^δ and other objects of QFT can be constructed. It is accepted [9] that in the limit $\delta \to 0$ this construction gives the solution of a quantization problem for the initial nonlocal system with the Lagrangian (13).

The first step is to represent the meromorphic function $\left[\mathcal{E}^\delta\right]^{-1}$ in the form

$$\left[\mathcal{E}^{\delta}(p^2)\right]^{-1} = \sum_{j=0}^{\infty}(-1)^j\frac{\mathcal{A}_j(\delta)}{m_j^2(\delta)-p^2} \tag{17}$$

where

$$m_j^2(\delta) = \frac{(j+1)^\sigma}{\delta}M_0^2 \tag{18}$$

and $\sigma < 1/\rho \le 2$, $M_0 = l_0^{-1}$. The coefficients $\mathcal{A}_j(\delta)$ can be easily calculated for any concrete function. Define the infinite set of Pais – Uhlenbeck auxiliary fields

$$\Phi_j^\delta(x) = \sqrt{\mathcal{A}_j(\delta)}\,\frac{\mathcal{E}^\delta\left(D^2\right)}{D^2+m_j^2(\delta)}\,\Phi^\delta(x) \tag{19}$$

$$\Phi^\delta(x) = \sum_{j=0}^{\infty}(-1)^j\sqrt{\mathcal{A}_j(\delta)}\,\Phi_j^\delta(x) \tag{20}$$

In terms of the fields Φ_i^δ the regularized Lagrangian (13) reads

$$\begin{aligned}-\mathcal{L}^\delta &= \sum_{j=0}^{\infty}(-1)^j\,\Phi_j^{\delta\dagger}(x)\left(D^2 + m_j^2(\delta)\right)\Phi_j^\delta(x) - \\ &\quad - \omega^\delta\Phi^{\delta\dagger}(x)*\Phi^\delta(x) + \lambda\left(\Phi^{\delta\dagger}(x)*\Phi^\delta(x)\right)^2\end{aligned} \tag{21}$$

Considering Φ_i^δ to be independent fields we can quantize the theory with this Lagrangian in the framework of canonical formalism. Then the equal-time canonical commutation relations take the form

$$\left[\Phi_i^{\delta\dagger}(0,\vec{x}),\ \Phi_j^\delta(0,\vec{y})\right] = i(-1)^j\delta_{ij}\delta\left(\vec{x}-\vec{y}\right) \tag{22}$$

Due to the presence of the factor $(-1)^j$ these relations can be realized only on the Hilbert space with the negative-norm quantum states. Moreover, it can be shown [9] that on this space the quantum field Φ^δ is a local operator.

Considering the SBGS we make standard shift of all fields Φ_j^δ independently

$$\Phi_j^\delta(x) = \left(\eta_j^\delta(x) + v_j^\delta\right)\frac{\chi}{\sqrt{2}} \tag{23}$$

$$\Phi^\delta(x) = \left(\eta^\delta(x) + v^\delta\right)\frac{\chi}{\sqrt{2}} \tag{24}$$

$$\chi = \begin{pmatrix}0\\1\end{pmatrix}$$

where

$$\eta^\delta(x) = \sum_{j=0}^{\infty}(-1)^j\sqrt{\mathcal{A}_j(\delta)}\,\eta_j^\delta(x) \tag{25}$$

$$v^\delta(x) = \sum_{j=0}^{\infty}(-1)^j\sqrt{\mathcal{A}_j(\delta)}\,v_j^\delta(x) \tag{26}$$

Redefining the theory in terms of the fields η_j^δ we fix the values of v_i^δ and ω^δ parameters from the requirement of the absence in the resulting

Lagrangian of any terms which are linear or quadratic in fields η_j^δ except $m_i^2(\delta)\eta_i^{\delta 2}$. Then we obtain the equations

$$m_i^2(\delta)v_i^\delta - v^\delta\sqrt{A_j(\delta)}\left(\omega^\delta - \lambda\left(v^\delta\right)^2\right) = 0 \tag{27}$$
$$\omega^\delta = 3\lambda\left(v^\delta\right)^2 \tag{28}$$

The solution of these equations reads

$$\omega^\delta = \frac{3}{2}\mathcal{E}^\delta(0) \tag{29}$$
$$\left(v^\delta\right)^2 = \frac{\mathcal{E}^\delta(0)}{2\lambda} \tag{30}$$
$$\left(v_i^\delta\right)^2 = \frac{A_j(\delta)\left(\mathcal{E}^\delta(0)\right)^3}{2\lambda m_i^4(\delta)} \tag{31}$$

Taking this into account and substituting (25) - (26) to (23) one finds after some transformations the final Lagrangian in the unitary physical gauge

$$\begin{aligned}
-\quad \mathcal{L}^\delta &= \frac{1}{2}\sum_{j=0}^{\infty}(-1)^j\,\eta_j^\delta\left(\partial^2 + m_j^2(\delta)\right)\eta_j^\delta + \lambda v^\delta\left(\eta^\delta\right)^3 + \frac{\lambda}{4}\left(\eta^\delta\right)^4 - \\
&- \frac{g^2}{4}\sum_{j=0}^{\infty}(-1)^j\left[\left(\eta_j^\delta\right)^2 + 2\eta_j^\delta v_j^\delta\right]\left(W_\mu^+W^{-\mu} + \frac{1}{2\cos^2\theta_W}Z_\mu Z^\mu\right) - \\
&- M_W^2\,W_\mu^+W^{-\mu} - \frac{1}{2}M_Z^2\,Z_\mu Z^\mu - \\
&- \sum_f \bar{\psi}_f\left(i\gamma_\mu D^\mu - m_f\left(1+\frac{\eta^\delta}{v^\delta}\right)\right)\psi_f
\end{aligned} \tag{32}$$

For the completeness the fermion part of the model is also included in this expression. The masses of the W and Z – bosons are given by

$$M_W^2 = \frac{g^2v^2\,\mathcal{E}'(0)}{4} = \frac{g^2\,\mathcal{E}(0)\mathcal{E}'(0)}{8\lambda} \tag{33}$$
$$M_Z^2 = \frac{M_W^2}{\cos^2\theta_W} \tag{34}$$

The Lagrangian (32) gives the solution of the problem we formulated at the beginning. In the limit $\delta \to 0$ it describes the SM with a nonlocal Higgs sector. The scalar particles corresponding to the quantized fields $\eta_j^\delta(x)$ escape from the observable spectrum for their masses $m_j^2(\delta) \to \infty$ when $\delta \to 0$. Moreover, the kinetic terms of these fields in the Lagrangian (32) can be written in the form

$$\sum_{j=0}^{\infty}(-1)^j\,\eta_j^\delta(x)\left(\partial^2 + m_j^2(\delta)\right)\eta_j^\delta(x) = \eta^\delta(x)\mathcal{E}^\delta\left(\partial^2\right)\eta^\delta(x) \tag{35}$$

Thus in the limit $\lim_{\delta\to 0}\eta^\delta = \eta$ one obtains the field η with the corresponding propagator $\mathcal{E}^{-1}(p^2)$ is being an entire function. This is another condition of absence of the observable scalar particles. From formulae (32), (34) it can be deduced that our model differs from the conventional SM only within the Higgs sector and its interactions with the gauge fields. The mass formulae for the W and $Z-$ bosons slightly differ but this difference is not essential as it will be shown further.

The calculations of the S-matrix elements in our model mainly are based on the standard technique of QFT perturbation theory.

The Feynman rules can be extracted directly from the Lagrangian (32). Here we write down only the $\eta_j^\delta(x)$ – field propagators:

$$\begin{aligned}
G^\delta_{ij(c)}(x-y) &= i < 0|\,T\left(\eta_i^\delta(x)\,\eta_j^\delta(y)\right)|0> = \\
&= (-1)^j\,\delta_{ij}\,\mathcal{D}^\delta_{j(c)}(x-y)
\end{aligned} \tag{36}$$

where

$$\mathcal{D}^\delta_{jc}(x) = \int\frac{d^4k}{(2\pi)^4}\frac{e^{-ikx}}{m_j^2(\delta) - k^2 - i\epsilon} \tag{37}$$

is the propagator of the local scalar field of the mass $m_j^2(\delta)$.

The essential peculiarities of the calculations in the framework of our nonlocal modification of the minimal SM are as follows. The diagrams with the external η_j^δ-lines must be excluded, but all possible internal

η_j^δ-states must be summed up. Final formulae for the physical matrix elements are obtained in the limit $\delta \to 0$.

It is relevant to note that the diagrams with the internal η_j^δ – lines are less divergent in comparison with the conventional SM. In some cases these diagrams are found to be ultraviolet finite. This is because of the fast decrease of the Euclidean Green functions in the NLQFT. In the present paper this property is expressed by the condition (6). If all the interactions are supposed to be nonlocal as accepted in the above-mentioned approach of Moffat [7], then the theory becomes finite.

3 Model Parameters and Domain of Perturbative Regime

To investigate the main features of the proposed nonlocal modification of the SM let us consider the simplest "minimal" variant of nonlocality. It corresponds to the following choice $\mathcal{W}(z) = z/M_0^2$ in formula (7). This gives

$$\mathcal{E}(z) = \mu^2 \exp(z/M_0^2) \tag{38}$$

where μ and M_0 are free parameters, extracted from the experiment. Regarding their possible values one notes that the μ-parameter has no essential physical meaning. The proper η-fields and the λ-coupling constant redefinition change it to any arbitrary value. In fact, the Lagrangian (32) is invariant under the following transformation

$$\eta^\delta \longrightarrow \eta'^\delta = \kappa^{-\frac{1}{2}}\,\eta^\delta \tag{39}$$

$$\eta_i^\delta \longrightarrow \eta'^\delta_i = \eta_i^\delta \tag{40}$$

$$\lambda \longrightarrow \lambda' = \kappa^2\,\lambda \tag{41}$$

$$\mathcal{E}^\delta(\partial^2) \longrightarrow \mathcal{E}'^\delta(\partial^2) = \kappa\,\mathcal{E}^\delta(\partial^2) \tag{42}$$

The latter transformation results in $\mu^2 \longrightarrow \mu'^2 = \kappa\mu^2$. In the considered modification of the SM there are no asymptotic conditions on the normalization of free fields for the absence of observable scalar particles. Therefore the physics is independent of the scale factor κ and value of μ^2. The only role of the μ – parameter is to provide a correct dimension of the propagator of the scalar field η. For the simplicity we take $\mu = M_0$. Then the mass formulae (33)– (34) acquire the same form as in the conventional SM.

For the practical calculations in the framework of our approach it is important to know the limitations on applicability of perturbation theory. These limitations come from the partial wave unitarity [12], [1] which require

$$|a_J(s)| \leq 1 \tag{43}$$

To find these limitations it is enough in our case to calculate the zeroth partial wave amplitude a_0 for longitudinal W and Z – boson scattering $W_L W_L \longrightarrow Z_L Z_L$. For the Lagrangian (32) with the $\mathcal{E}$-function in the form (38) we have in tree approximation

$$a_0 = \frac{g^2}{64\pi}\,\frac{M_0^2}{M_W^2}\left(e^{S/M_0^2} - 1\right) \tag{44}$$

this gives the limitation

$$S \leq S_{max}(M_0) = M_0^2\, ln\left(\frac{8\pi\sqrt{2}}{G_F M_0^2} + 1\right) \tag{45}$$

In this kinematic domain the condition $|a_0(s)| \leq 1$ is satisfied. The function $S_{max}(M_0)$ monotonically increases with M_0 and asymptotically tends to a constant

$$\lim_{M_0\to\infty} S_{max}(M_0) = \frac{8\sqrt{2}\pi}{G_F} \simeq (1.8\,TeV)^2 \tag{46}$$

The saturation occurs rather quickly, and for $M_0 = 1.5$ TeV we get the upper bound $\sqrt{S} \leq \sqrt{S_{max}} \simeq 1.5$ TeV. To achieve higher energies it is necessary to perform the calculations in the next to leading order of perturbation theory and maybe to sum up some classes of diagrams.

Note that the obtained limitation (45) strongly depends on the form of the $\mathcal{E}$-function. Starting from the function differing from (38) we arrive at the limitation other than (45).

In the conventional local SM the unitarity gives an upper bound on the Higgs boson mass. At the tree level this is $M_H \leq 2$ TeV. If the

Higgs boson is not discovered in this mass region, then perturbation theory is not applicable for the derivation of the SM predictions at high energy. Moreover, some of the low energy predictions obtained within this framework are subject to doubt. The reason is as follows. Because in the CM there is a well known relation

$$M_H^2 = 2\lambda v_{SM}^2 \tag{47}$$

where $v_{SM} \simeq 250$ GeV, then for $M_H > 1$ TeV we have $\frac{\lambda}{4\pi^2} > 1$. The last quantity characterizes perturbative corrections due to the Higgs self-interaction. Therefore, if it is not small enough, the perturbation theory fails. This conclusion does not depend on the kinematic domain considered but only on the M_H – value. As we have seen, in the proposed nonlocal modification of SM the situation is different. There is no connection between Higgs invisibility and applicability of the perturbation theory. It has been shown that the lowest order calculations are correct in the definite kinematic domain which depends on the form of nonlocality in the Higgs self-interaction. For the minimal variant (38) the domain is defined by (45) .

4 Conclusion

Thus, the introduction of nonlocality in the self-interaction of Higgs fields enables us to exclude the scalar Higgs particles from the observable spectrum of the SM. We don't specify the nature of this nonlocality. In principle, it may be considered not a fundamental physical notion but an effective phenomenological way of taking into account some of interactions beyond the SM . Proposed nonlocal modification of the SM is described by the renormalizable Lagrangian (32) for which we have formulated the rules of perturbative calculations of physical matrix elements. Though the applicability of the lowest order calculations is bounded within the energy domain (45) in this framework one can, nevertheless, calculate most of the electroweak effects and compare the predictions with the experimental data. These will be considered elsewhere in our next paper. Here we confine ourselves to the following general remarks. The main physical difference between the proposed nonlocal modification of the SM and the conventional one comes from their Higgs sectors. However, the Higgs field interacts with the leptons and quarks, except the heavy t-quark, rather weakly. Thus we may expect the observable difference in predictions of these two variants of the SM at the level of refined effects of the radiative corrections or in the processes difficult for the experimental investigation. Among them there are widely discussed processes of the W and Z – boson scattering. They will be accessible in the near future at the SSC and LHC – colliders where it is planned to search for possible growth with the energy of their cross sections. This behaviour is predicted in the conventional SM for the longitudinal vector boson scattering, for example, $W_L W_L \longrightarrow Z_L Z_L$. While the nonlocal modification predicts the growth for the every polarizations of W and Z – bosons. Another discrepancy is the radiative corrections. They also depend on the nature of the Higgs sector. One can expect the most considerable difference in the predictions of the nonlocal modification and the conventional SM for those radiative effects which, being calculated in the conventional SM, have the strongest dependence on the Higgs boson mass.

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