Dedicated to our friend
Nikita Sveshnikov

Composite Two-Higgs Models with Flavour Changing Neutral Currents Suppression

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1. Introduction

The mass generation in the Standard Model (SM) of electroweak interactions is provided by the Higgs mechanism which rests on the Electroweak Symmetry Breaking (EWSB). To accommodate the well-established electromagnetic and weak phenomena, the Higgs mechanism requires the existence of at least one iso-doublet, complex scalar field (Higgs boson). If SM holds valid as a weak-coupling theory till very high energies then the massive component of this iso-doublet is a neutral scalar particle which cannot have a mass heavier than few hundreds of GeV. Thus the search for the Higgs boson is one of the fundamental quests for testing the minimal SM. Current estimations based on the different theoretical requirements and experimental implications[1, 2] give the SM Higgs mass in the "intermediate mass" window $65 < M_H < 200\,GeV$ for a top quark mass value of about $175\,GeV$ [3].

Despite of the recent successes of the SM in its excellent agreement with the precision measurements at present energies [4], it is generally believed that the SM is not the final theory of elementary particle interactions. There are many extensions of the SM which lead to the enlargement of the Higgs sector of the SM. For instance, the Minimal Supersymmetric Standard Model (MSSM) [5] entails two elementary Higgs doublets at low energies, the Two-Higgs-Doublet Model (2HDM) contains two complex $SU(2)_L$ - doublet scalar fields with hypercharge $Y = \pm 1$ to couple the up-type / down-type right-handed quarks to its Higgs doublet. The search for relations between the many Higgs-field dynamics and the masses of t-quark and Higgs boson give the selection rule for a particular model beyond the SM as well as for its acceptable parameters [6], [7], [8], [9]. In more complicated theories (see [1] and references therein) such as SUSY SM ones, E_6 ones, or Left-Right symmetric ones [10], several neutral scalars, charged scalars and even double-charged scalars are required in order to give all amplitudes acceptable high-energy behavior [1, 2].

However there exists an alternative possibility [11] to restrict the number of elementary particles to the observable fermion and vector-boson sector with generation of composite scalar Higgs particles due to attractive self-fermion interaction. The idea that the Higgs boson could be a bound state of heavy quark pairs has been developed and worked out in a series of papers [11], being motivated by the earlier work of Nambu and Jona- Lasinio (NJL) [12]. In particular, for t- quarks, it is provided by the Top-Mode Standard Model (TSM) Lagrangian, known also as the Bardeen-Hill-Lindner (BHL) Lagrangian [11],[13],[14]. The possibility that multiple four-fermion interactions (for three and a heavy fourth generations) are important in EWDSB, leading to an effective 2HDM at low energies, has been investigated in [15]. In this model Higgs boson induced Flavour Changing Neutral Currents (FCNS's) are naturally suppressed [16]. Some recent theoretical aspects and questions of $\bar{t}t$- condensation frameworks

one can find in the review of [17]. In these scenarios the heavy top mass is explained by the "top-condensation" where new strong forces lead to the formation of $\bar{t}t$ bound states and the EWSB. In a minimal version of quark models the top-condensation was triggered by a local four-fermion interaction.

We propose the quark models with *Quasilocal* four-fermion interaction [18] where the derivatives of fermion fields are included into vertices to influence on the formation of the second Higgs doublet. Such extensions of the Higgs sector lead to a broad spectrum of excited bound states, moreover they may be viewed as more natural than other, above mentioned extensions since the particles involved in EWSB form only a ground state spectrum generic for SM. In these quasilocal NJL-like quark models (QNJLM) the symmetries do not forbid further higher dimensional vertices[18, 20] and one should expect that the ground states could be accompanied by (radial) excitations with identical quantum numbers but much higher masses [19, 21, 22].

Thus, from the viewpoint of the 2HDM SM, the QNJLM are attractive because: i) it is an extension of the minimal TSM which adds new phenomena (e.g. a broad mass spectrum of bound states including charged Higgs bosons); ii) it is a minimal extension in that it adds the fewest new arbitrary constants; iii) it easily satisfies theoretical constraints on $\rho \simeq 1$ and the tree-level FCNC's suppression [16] in accordance with the experimental evidence; iv) such a Higgs structure is required in order to build a model with the CP- violation [23] because the one-Higgs doublet interaction does not provide any effect of dynamical CP- violation. We shall show in a toy model with quasilocal four-fermion interaction how P-parity breaks down dynamically for the special choice of coupling constants [24].

In our talk we give the description of the design of the Quasilocal Quark Models of type I and type II which provide two composite Higgs doublets and in principle satisfy phenomenological restrictions on FCNC suppression.

2. Quasilocal Quark Models and Polycritical Regime

In order to involve in the theory the effects of the discarded states at scales of order Λ it is needed to adjust the existing couplings constants in the Lagrangian and to add new, quasilocal, non-renormalizable interactions (vertices). These vertices are polynomial in the fields and derivatives of the fields and only a finite number of interactions is required when working to a particular order in $\frac{p}{\Lambda}$, where p is a typical momentum in whatever process is under study.

We examine the DCSB patterns in the mean-field approach (large-N_c limit) and estimate the vertices with any number of fermion legs and derivatives. The main rule to select out relevant vertices is derived from the requirement of insensitivity in respect to the separation scale Λ following the conception of low-energy effective action [18].

We assume that:

(i) Λ^2-order contributions from different vertices are dominant in creating the DCSB-critical surface that is provided by cancellation of all contributions of Λ^2-order and defines the polycritical regime;

(ii) Λ^0-order contributions from vertices assemble in the mean-field action to supply fermions with dynamical mass $m_d << \Lambda$ which establishes the low-energy physical scale;

(iii) respectively Λ^{-2} (etc.)-order contributions are irrelevant at energies much lower than Λ and so may be dropped from the theory if such accuracy is unnecessary.

In the large-N_c approach the following approximation for v.e.v. of fermion operators is valid,

$$\langle(\bar{q}q)^n\rangle \;=\; \Big(\langle\bar{q}q\rangle\Big)^n \Big(1 \;+\; O(1/N_c)\Big), \tag{1}$$

where any number of derivatives can be inserted between antifermion and fermion operators.

V.e.v. of a bilinear operator is estimated in the assumption that quarks obtain a dynamical mass. Namely,

$$\langle\bar{q}\left(\frac{\partial^2}{\Lambda^2}\right)^n q\rangle \;\sim\; \frac{1}{\Lambda^{2n}} \int\limits_{|p|<\Lambda} \frac{d^4p}{(2\pi)^4}\; tr\frac{p^{2n}}{\not{p} + im_d} \;\sim\; N_c m_d \Lambda^2. \tag{2}$$

One can see that the vertices with derivatives in many-fermion interaction are not suppressed and play equal role in the mass-gap equation.

We omit the full classification of effective vertices (see [18]) and describe only the minimal structure of the QNJLM which admits the polycritical regime,

$$\mathcal{L} = \bar{q}\,\not{D} q + \frac{1}{N_c \Lambda^2} \sum_{m,n=0}^{l} a_{mn}\, \bar{q}_R f_n \left(\frac{-\partial^2}{\Lambda^2}\right) q_L \cdot \bar{q}_L f_m \left(\frac{-\partial^2}{\Lambda^2}\right) q_R, \tag{3}$$

where $\partial \equiv 1/2\,(\overrightarrow{\partial} - \overleftarrow{\partial})$ and a_{mn} is a hermitian matrix of coupling constants without zero eigenvalues. It is taken to be real symmetric one in order that the interaction did not break the CP-parity explicitely. Chiral fermion fields are given by $q_{L(R)} = 1/2(1 \pm \gamma_5)q$. We define the vertex formfactors to be polynomials of derivatives,

$$f_m(\tau) = \sum_{i=0}^{K_m} f_m^{(i)} \tau^i, \tag{4}$$

to have quasilocal interactions. The variable τ is related to derivatives, $\tau \to -\partial^2/\Lambda^2$.

Let us regularize the interaction vertices with the help of a momentum cutoff, $\bar{q}q \longrightarrow \bar{q}\theta(\Lambda^2 + \partial^2)q$. Without loss of generality one can choose formfactors $f_i(\tau)$ being orthogonal polynomials on the unit interval,

$$\int_0^1 d\tau f_m(\tau) f_n(\tau) = \delta_{mn}. \tag{5}$$

Let us now introduce the appropriate set of auxiliary fields $\phi_n(x) \sim const$ and develop the mean-field approach,

$$\mathcal{L}(\phi) \;= \bar{q}\Big(\not{D} + iM(\phi)P_L + iM^+(\phi)P_R\Big)q + N_c\Lambda^2 \sum_{m,n=1}^{l} \phi_m^*\, a_{mn}^{-1}\, \phi_n\,. \tag{6}$$

The dynamical mass functional is a linear combination of formfactors,

$$M(\phi) \equiv \sum_{n=1}^{l} \phi_n(x) f_n\left(\frac{-\partial^2}{\Lambda^2}\right). \tag{7}$$

Thereby we come to a model with l channels. When integrating out the fermion fields one obtains the effective action of ϕ^*, ϕ - fields. The effective potential V_{eff} is proved to be a functional depending on the dynamical mass functional $M(\phi^*, \phi)$ and proportional to N_c that allows us to use the saddle point approximation for $N_c >> 1$.

The effective potential for auxiliary fields can be derived with the momentum cutoff regularization by averaging over quark fields:

$$\begin{aligned} V_{eff}(\phi) &= N_c\Lambda^2 \sum_{m,n=1}^{l} \phi_m^\star a_{mn}^{-1} \phi_n - \frac{N_c\Lambda^4}{8\pi^2} \int_0^1 d\tau\tau \ln\left(1 + \frac{|M(\tau)|^2}{\Lambda^2\tau}\right) \\ &= \frac{N_c}{8\pi^2}\Bigg[\Lambda^2 \sum_{m,n=1}^{l} \phi_m^\star (8\pi^2 a_{mn}^{-1} - \delta_{mn})\phi_n + \frac{1}{2}|M_0|^4 \left(\ln\frac{\Lambda^2}{|M_0|^2} + \frac{1}{2}\right) \\ &\quad + \frac{1}{2}\int_0^1 \frac{d\tau}{\tau}(|M(\tau)|^4 - |M_0|^4) + O\left(\frac{1}{\Lambda^2}\right)\Bigg], \end{aligned} \tag{8}$$

herein $M_0 \equiv M(0)$. The last approximation is valid in such a strong coupling regime where the dynamical mass $M_0 << \Lambda$. This regime is of our main interest and it is realized in the vicinity of a (poly)critical surface. The critical values of coupling constants, $a_{mn}^c = \delta_{mn}/8\pi^2$, are found from the cancellation of

quadratic divergences. Here we study the critical regime in all l channels. The vicinity of this polycritical point is described by the following parametrization:

$$8\pi^2 a^{-1}_{mn} \sim \delta_{mn} + \frac{\Delta_{mn}}{\Lambda^2}, \qquad |\Delta_{ij}| \ll \Lambda^2. \tag{9}$$

The generalized mass gap equations,

$$\frac{\delta V_{eff}(\phi,\phi^*)}{\delta \phi^*_m} = 0 = \frac{\delta V_{eff}(\phi,\phi^*)}{\delta \phi_m}, \tag{10}$$

deliver the extremum to the effective potential which may cause the DCSB if it is an absolute minimum. They read:

$$\begin{aligned}
\sum_{n=1}^{l} \Delta_{mn}\phi_n &= \int_0^1 \frac{d\tau}{\tau + \frac{|M(\tau)|^2}{\Lambda^2}} (|M(\tau)|^2 M(\tau) f_m(\tau) \\
&\simeq f_m(0)|M_0|^2 M_0 \ln\frac{\Lambda^2}{M_0^2} \\
&+ \int_0^1 \frac{d\tau}{\tau}(|M(\tau)|^2 M(\tau) f_m(\tau) - |M_0|^2 M_0 f_m(0)).
\end{aligned} \tag{11}$$

The true minimum is derived from the positivity of the second variation of the effective action around a solution of the mass-gap equation,

$$\phi_m = <\phi_m> + \sigma_m(x) + i\pi_m(x). \tag{12}$$

This variation reads:

$$\begin{aligned}
\frac{4\pi^2}{N_c}\delta^2 S_{eff} \quad &\equiv \left(\sigma, (A^{\sigma\sigma}p^2 + B^{\sigma\sigma})\sigma\right) \\
&+2\left(\pi, (A^{\pi\sigma}p^2 + B^{\pi\sigma})\sigma\right) + \left(\pi, (A^{\pi\pi}p^2 + B^{\pi\pi})\pi\right),
\end{aligned} \tag{13}$$

where two symmetric matrices - for the kinetic term $\hat{A} = \left(A^{ij}_{mn}\right), \quad i,j = (\sigma,\pi)$ and for the constant, momentum independent part, $\hat{B} = \left(B^{ij}_{mn}\right)$ - have been introduced.

The positivity of the second variation corresponds to the formation of physical mass spectrum for composite scalar and pseudoscalar states which can be found from zeroes of the second variation determinant at the Minkovski momenta ($p^2 < 0$),

$$\det(\hat{A}p^2 + \hat{B}) = 0. \tag{14}$$

Matrix elements of $\hat{B}$ are given by the following relations:

$$\begin{aligned}
B^{\sigma\sigma}_{mn} &= 6\int_0^1 \frac{d\tau}{\tau}\left[(\,\mathrm{Re}M)^2 f_m(\tau) f_n(\tau) - M_0^2 f_m(0) f_n(0)\right] \\
&+ M_0^2 f_m(0) f_n(0)\left(6\ln\frac{\Lambda^2}{M_0^2} - 4\right) - 2\Delta_{mn} \\
&+ 2\int_0^1 \frac{d\tau}{\tau}(\,\mathrm{Im}M)^2 f_m(\tau) f_n(\tau),
\end{aligned} \tag{15}$$

$$\begin{aligned}
B^{\pi\pi}_{mn} &= 2\int_0^1 \frac{d\tau}{\tau}\left[(\,\mathrm{Re}M)^2 f_m(\tau) f_n(\tau) - M_0^2 f_m(0) f_n(0)\right] \\
&+ 2M_0^2 f_m(0) f_n(0)\ln\frac{\Lambda^2}{M_0^2} - 2\Delta_{mn} \\
&+ 6\int_0^1 \frac{d\tau}{\tau}(\,\mathrm{Im}M)^2 f_m(\tau) f_n(\tau),
\end{aligned} \tag{16}$$

$$B^{\sigma\pi}_{mn} = 4\int\limits_0^1 \frac{d\tau}{\tau}(\mathrm{Re}M)(\ \mathrm{Im}M) f_m(\tau) f_n(\tau), \tag{17}$$

where the terms of $1/\Lambda^2$-order are neglected.

When exploiting the mass-gap equation (11) one can prove that the matrix $\hat{B}$ has always a zero eigenvalue related to the eigenvector $\phi^0_m =< \pi_m > -i\cdot < \sigma_m >$. It corresponds to the arising of the Goldstone mode (the massless Goldstone bosons).

One can convince oneself that if the solution $M(\tau)$ has an imaginary part the mixing between scalar and pseudoscalar channels arise which leads to the P-parity breaking. Of course, this "axial" mass should not be removed by a global chiral rotation.

The kinetic energy matrix $\hat{A}$ turns out to be block-diagonal [24], and does not induce any P-parity breaking.

$$A^{\sigma\pi}_{mn} = 0, \quad A^{\pi\pi}_{mn} = A^{\sigma\sigma}_{mn}\left[1 + O\left(\frac{1}{\ln\frac{\Lambda^2}{\mu^2}}\right)\right], \tag{18}$$

$$\begin{aligned} A^{\sigma\sigma}_{mn} &= \frac{1}{2}\Bigg[f_m(0)f_n(0)\left(\ln\frac{\Lambda^2}{M_0^2} + O\left(1\right)\right) \\ &+ \int\limits_0^1 \left[f_m(\tau)f_n(\tau) - f_m(0)f_n(0)\right]\frac{d\tau}{\tau}\Bigg] + O\left(\frac{1}{\Lambda^2}\right), \end{aligned} \tag{19}$$

herein we have displayed the leading terms only in the large-log approximation. The more detailed expression can be found in [24], [19].

3. Higgs Bosons as Radial Excitations - Model I

Let us construct now the two-flavor quark models with quasilocal interaction in which the t- and b-quarks are involved in the DCSB. In accordance with the SM, the left components of both quarks form a doublet:

$$q_L = \begin{pmatrix} t_L \\ b_L \end{pmatrix}, \tag{20}$$

which transforms under $SU(2)_L$ group as a fundamental representation while the right components t_R, b_R are singlets.

The Model I which is to satisfy the FCNC suppression has the following Lagrangian:

$$\begin{aligned} \mathcal{L}_{\mathcal{J}} &= \bar{q}_L \not{p} q_L + \bar{t}_R \not{p} t_R + \bar{b}_R \not{p} b_R + \\ &\frac{8\pi^2}{N_c\Lambda^2}\sum_{k,l=1}^{2} a_{kl}\left(g_{t,k}J^T_{t,k} + g_{b,k}\widetilde{J}^T_{b,k}\right) i\tau_2\left(g_{b,l}J_{b,l} - g_{t,l}\widetilde{J}_{t,l}\right). \end{aligned} \tag{21}$$

Here we have introduced the denotations for doublets of fermion currents:

$$J_{t,k} \equiv \bar{t}_R f_{t,k}\left(-\frac{\partial^2}{\Lambda^2}\right) q_L, \qquad J_{b,k} \equiv \bar{b}_R f_{b,k}\left(-\frac{\partial^2}{\Lambda^2}\right) q_L, \tag{22}$$

and the tilde in $\widetilde{J}_{t,k}$ and $\widetilde{J}_{b,k}$ marks charge conjugated quark currents, roteted with τ_2 Pauli matrix

$$\tilde{J}_{t,k} = i\tau_2 J^\star_{t,k}, \qquad \widetilde{J}_{b,k} = i\tau_2 J^\star_{b,k} \tag{23}$$

The subscripts t, b indicate right components of t and b quarks in the currents, the index k enumerates the formfactors:

$$\begin{aligned} f_{t,1} &= 2 - 3\left(-\frac{\partial^2}{\Lambda^2}\right), \quad f_{t,2} = -\sqrt{3}\left(-\frac{\partial^2}{\Lambda^2}\right), \\ f_{b,1} &= 2 - 3\left(-\frac{\partial^2}{\Lambda^2}\right), \quad f_{b,2} = -\sqrt{3}\left(-\frac{\partial^2}{\Lambda^2}\right). \end{aligned} \tag{24}$$

As the spinor indices are contracted to each other in (22), $J_{t,k}$ transforms as a doublet under $SU(2)_L$. τ_2 is a Pauli matrix in the adjoint representation of the group $SU(2)_L$. Coupling constants of the four-fermion interaction are represented by 2×2 matrix a_{kl} and contributed also from the Yukawa constants $g_{k,l}, g_{b,k}$.

The Lagrangian density of the Model I (21) to describe the dynamics of composite Higgs bosons can be obtained by means of introduction of auxiliary bosonic variables and by integrating out fermionic degrees of freedom. According to this scheme, we define two scalar $SU(2)_L$-isodoublets:

$$\Phi_1 = \begin{pmatrix} \phi_{11} \\ \phi_{12} \end{pmatrix}, \qquad \Phi_2 = \begin{pmatrix} \phi_{21} \\ \phi_{22} \end{pmatrix} \tag{25}$$

and their charge conjugates:

$$\widetilde{\Phi}_1 = \begin{pmatrix} \phi^\star_{12} \\ -\phi^\star_{11} \end{pmatrix}, \qquad \widetilde{\Phi}_2 = \begin{pmatrix} \phi^\star_{22} \\ -\phi^\star_{21} \end{pmatrix}. \tag{26}$$

In terms of auxiliary fields, the Lagrangian (21) can be rewritten in the following way:

$$\mathcal{L}_{\mathcal{J}} = L_{kin} + \frac{N_c \Lambda^2}{8\pi^2} \sum_{k,l=1}^{2} \Phi^\dagger_k (a^{-1})_{kl} \Phi_l + i \sum_{k=1}^{2} \Big[g_{t,k} \widetilde{\Phi}^\dagger_k J_{t,k} + g_{b,k} \Phi^\dagger_k J_{b,k} \Big] + h.c. \tag{27}$$

The integrating out of fermionic degrees of freedom will produce the effective action for Higgs bosons of which we shall keep only the kinetic term and the effective potential consisting of two- and four-particles vertices. The omitted terms are supposedly small, being proportional to inverse powers of a large scale factor Λ. The Yukawa constants are chosen of the form

$$g_{t,k} = 1; \qquad g_{b,k} = g \tag{28}$$

for $k = (1, 2)$. The first constant is set to unit because the fields Φ_1 and Φ_2 can always be multiplied by an arbitrary factor which is absorbed by four fermion coupling constants through redefinition of polycritical coupling constants. The remaining constant g induces thereby the quark mass ratio m_b/m_t.

We assume the electric charge stability of vacuum or, in other words, that only neutral components of both Higgs doublets may have nonzero v.e.v. Hence, one can deal with only neutral components of the Higgs doublets in the effective action for studying DCSB. This part of the Higgs sector can be investigated separately as a model where two singlets (not doublets) appear as composite Higgs bosons. For this purpose, we use the Quasilocal Two-Channel model[19].

Following the definitions made in [24], we relate the fields ϕ_1, ϕ_2 and ρ to the neutral components of Higgs doublets:

$$\phi_{11} = \phi_1; \qquad \phi_{22} = \phi_2 + i\rho. \tag{29}$$

The condition of minimum of the effective potential with the charged components of Higgs doublets put to zero values: $\phi_{12} = \phi_{21} = 0$, brings the mass-gap equations for them:

$$\begin{aligned}
\Delta_{11}\phi_1 + \Delta_{12}\phi_2 &= 16\phi_1^3 \ln\frac{\Lambda^2}{4\phi_1^2} + 16g^4\phi_1^3 \ln\frac{\Lambda^2}{4g^2\phi_1^2} - (1+g^4)\Big[\frac{159}{4}\phi_1^3 - \frac{15\sqrt{3}}{4}\phi_1^2\phi_2 + \\
&\quad + \frac{9}{4}\phi_1\phi_2^2 + \frac{\sqrt{3}}{4}\phi_2^3 + \frac{\rho^2}{4}(3\phi_1 + \sqrt{3}\phi_2)\Big], \qquad (30) \\
\Delta_{12}\phi_1 + \Delta_{22}\phi_2 &= (1+g^4)\Big[-\frac{5\sqrt{3}}{4}\phi_1^3 + \frac{9}{4}\phi_1^2\phi_2 + \frac{3\sqrt{3}}{4}\phi_1\phi_2^2 + \\
&\quad + \frac{9}{4}\phi_2^3 + \frac{\rho^2}{4}(\sqrt{3}\phi_1 + 9\phi_2)\Big], \qquad (31) \\
4\rho\Delta_{22} &= \rho(1+g^4)\left(3\phi_1^2 + 2\sqrt{3}\phi_1\phi_2 + 9\phi_2^2 + 9\rho^2\right) \qquad (32)
\end{aligned}$$

Let us consider the equations (30)–(32) for two cases: 1) $\rho = 0$ and 2) $\rho \neq 0$.

When $\rho = 0$, assuming that $\phi_1 \neq 0$ and $\phi_2 \neq 0$ one finds the following solutions.
The solution of the mass-gap equation of Gross-Neveu-type is:

$$\phi_1^2 \approx \frac{\det \Delta}{16(1+g^4)\Delta_{22} \ln\left(\frac{\Lambda^2}{\mu^2}\right)}, \quad \phi_2 \approx -\frac{\Delta_{12}}{\Delta_{22}}\phi_1. \tag{33}$$

The solution of the mass-gap equation of the Abnormal-type is:

$$\phi_2^2 \approx \frac{4}{9(1+g^4)}\Delta_{22}, \quad \phi_1^2 \approx \frac{\Delta_{22}^{1/3}(3\sqrt{3}\Delta_{12} - \Delta_{22})^{2/3}}{4 \cdot 3^{5/3}(1+g^4)\ln^{2/3}\left(\frac{\Lambda^2}{\mu^2}\right)}, \tag{34}$$

For the case 2, of a non-zero ρ , it can be shown that for fixed Δ_{kl} while Λ grows larger, the solution exists if Δ_{kl} parameters are chosen close to a particular plane in the parametric space. This plane is defined by the equation:

$$\Delta_{22} = 3\sqrt{3}\Delta_{12}. \tag{35}$$

When Δ_{kl} satisfy the equation (35) exactly, the solution is found to be as follows (in the large-log approximation):

$$\phi_1^2 = \frac{\Delta_{11}}{16(1+g^4)\ln\frac{\Lambda^2}{\mu^2}}\left[1 + O\left(\frac{1}{\ln\frac{\Lambda^2}{\mu^2}}\right)\right], \tag{36}$$

$$\phi_2 = \sqrt{3}\phi_1, \tag{37}$$

$$\rho^2 = \frac{4}{9}\frac{\Delta_{22}}{(1+g^4)}\left[1 + O\left(\frac{1}{\ln\frac{\Lambda^2}{\mu^2}}\right)\right]. \tag{38}$$

The mass spectrum of related bosonic states is determined by the Eqs. (26)-(31) and taking into account the conditions necessary for a minimum of the potential (52,53). The solutions at $-m^2 = p^2 < 0$ one can obtain from the Eqs:

$$\det(Ap^2 + B) = 0, \tag{39}$$

The "kinetic" matrix $\hat{A}$ as being proportional to p^2 is derived in the soft-momentum expansion in powers of p^2 and in a large-logΛ approximation. After substituting expressions for the matrix $\hat{A}, \hat{B}$ into (39) one can get the mass spectrum for the neutral Higgs bosons in Model I. For the case, $\rho = 0$, the mass spectrum is resembling ones in Two-Channel model[19], in particular, the Gross-Neveu-type solution brings the spectrum for scalars:

$$m_{\sigma'}^2 \approx -\frac{4\Delta_{22}}{3(1+g^2)}, \tag{40}$$

$$m_{\sigma}^2 \approx -\frac{\det \Delta}{(1+g^2)\Delta_{22}\ln\left(\frac{\Lambda^2}{\mu^2}\right)} = 4m_d^2 \tag{41}$$

and for pseudoscalars:

$$m_{\pi'}^2 \approx -\frac{4\Delta_{22}}{3(1+g^2)}, \tag{42}$$

$$m_{\pi}^2 = 0. \tag{43}$$

The Abnormal solution induces the following mass spectrum for scalars:

$$m_{\sigma'}^2 \approx \frac{8\Delta_{22}}{3(1+g^2)}, \tag{44}$$

$$m_{\sigma}^2 \approx \frac{2\Delta_{22}^{1/3}(3\sqrt{3}\Delta_{12} - \Delta_{22})^{2/3}}{3^{2/3}(1+g^2)\ln^{2/3}\left(\frac{\Lambda^2}{\mu^2}\right)} = 6m_d^2 \tag{45}$$

and for pseudoscalars:

$$m^2_{\pi'} \approx \frac{3^{2/3}(3\sqrt{3}\Delta_{12}-\Delta_{22})^{4/3}}{27(1+g^2)\Delta_{22}^{1/3}\ln^{1/3}\left(\frac{\Lambda^2}{\mu^2}\right)}, \quad (46)$$

$$m^2_{\pi} = 0 \quad (47)$$

Remark that dynamical mass m_d is in fact the mass of t quark in the Model-I, because the v.e.v. of ϕ_{11}, which is parametrized as $\phi_1 \equiv \langle\phi_{11}\rangle$, gives the value of mass of t-quark.

The mass spectrum in the P-parity Breaking Phase, for a non-zero ρ is:

$$m^2_1 = 0, \quad (48)$$

$$m^2_2 \approx \frac{3\Delta_{11}-\Delta_{22}}{3(1+g^2)\ln\frac{\Lambda^2}{\mu^2}} \approx 16\phi_1^2 = 4m_d^2, \quad (49)$$

$$m^2_3 \approx \frac{8\Delta_{22}}{3(1+g^2)}, \quad (50)$$

$$m^2_4 \approx \frac{3\Delta_{11}+7\Delta_{22}}{27(1+g^2)\ln\frac{\Lambda^2}{\mu^2}}, \quad (51)$$

Thus, we have constructed the Model I where:

a) Two composite Higgs doublets are created dynamically as a consequence of DCSB in two channels.

b) In 2HQ Model I Higgs bosons are rather radial, ground and excited states in the scalar-pseudoscalar channels.

c) The appropriate fine tuning leads also to spontaneous breaking of P-parity and, therefore, of CP-parity in the Higgs sector.

4. Top-Bottom Condensation for 2HQM Model- II

The Lagrangian density of the Model II to describe the dynamic of two composite Higgs bosons which consist of bound states (condensates $\bar{t}t, \bar{b}b$) and satisfy the FCNC suppression[16] can be written as:

$$\mathcal{L}_{\mathcal{J}} = L_{kin} + \frac{N_c\Lambda^2}{8\pi^2}\sum_{k,l=1}^{2}\Phi_k^\dagger(a^{-1})_{kl}\Phi_l + i\bar{q}(\widetilde{\mathbf{M}}P_L + \widetilde{\mathbf{M}}^\dagger P_R)q + h.c., \quad (52)$$

where $P_{L(R)} = 1/2(1\pm\gamma_5)$ - the left and right projectors, and $\widetilde{\mathbf{M}}$ is the two-by-two flavour matrix:

$$\widetilde{\mathbf{M}} = \sum_{m=1}^{2}\begin{pmatrix} g_{t,k}\phi_{m2}\, f_{t,m}\left(\frac{p^2}{\Lambda^2}\right) & -g_{t,k}\phi_{m1}\, f_{t,m}\left(\frac{p^2}{\Lambda^2}\right) \\ g_{b,k}\phi^\star_{m1}\, f_{b,m}\left(\frac{p^2}{\Lambda^2}\right) & g_{b,k}\phi^\star_{m2}\, f_{b,m}\left(\frac{p^2}{\Lambda^2}\right) \end{pmatrix}, \quad (53)$$

In this Model II Φ_1, Φ_2 give masses to $up-, down-$type quarks. The structure of quark interaction is specified in four formfactors:

$$f_{t,1} = 1 - c_{t,1}\frac{\partial^2}{\Lambda^2}, \qquad f_{t,2} = -c_{t,2}\frac{\partial^2}{\Lambda^2},$$
$$f_{b,1} = -c_{b,1}\frac{\partial^2}{\Lambda^2}, \qquad f_{b,2} = 1 - c_{b,2}\frac{\partial^2}{\Lambda^2}. \quad (54)$$

When the chiral symmetry is broken, the v.e.v. of neutral Higgs fields are non-zero and the true Yukawa vertices should be obtained by subtracting from $\widetilde{\mathbf{M}}$ its v.e.v.

$$\mathbf{M} = \widetilde{\mathbf{M}} - \langle\widetilde{\mathrm{M}}\rangle, \quad (55)$$

where $\langle \ldots \rangle$ means a v.e.v.:

$$\langle \widetilde{\mathbf{M}} \rangle = \begin{pmatrix} m_t & 0 \\ 0 & m_b \end{pmatrix}. \tag{56}$$

The elements of the matrix (56) are the quark mass functions:

$$\begin{aligned} m_t(\tau) &= \phi_1(1 + c_{t,1}\tau) + \phi_2 e^{i\delta_0} c_{t,2}\tau, \\ m_b(\tau) &= \phi_1 c_{b,1}\tau + \phi_2 e^{i\delta_0}(1 + c_{b,2}\tau), \qquad \tau \equiv -\frac{\partial^2}{\Lambda^2}, \end{aligned} \tag{57}$$

defined to be real and $\phi_1 =< \phi_{12} >, \phi_2 =< \phi_{22} >$. The non-zero phase at m_b, which is displayed explicitly in (56), may appear if the v.e.v. of ϕ_{22} acquires irremovable phase factor when the chiral symmetry is broken. As the vacuum charge stability is assumed, $\langle \widetilde{\mathrm{M}} \rangle$ is diagonal.

The effective potential of Model II and solutions of the corresponding mass-gap equations are obtained in the scheme described above for Model I. (For more details the reader is referred to [26]).

The suppression of the flavor changing neutral currents in a composite two-Higgs model can be implemented when Yukawa couplings are chosen as follows:

$$g_{t,1} = g_{b,2}; \qquad g_{t,2} = g_{b,1} = \sqrt{5}. \tag{58}$$

To be short we display only the mass spectrum. When the mass ratio of t- and b-quarks m_t/m_b holds unchanged while Λ goes to infinity, in the large-log approach we get the following estimations for Higgs boson masses:

1. for the case of conserved P-parity ($\delta_0 = 0$):

$$\begin{aligned} m_\sigma &= 2m_b, \\ m_{\sigma'} &\approx 2m_t, \\ m_\pi &= 0, \\ m_{\pi'} &= \sqrt{\frac{(m_t^2 + m_b^2)}{2m_t m_b \ln\frac{\Lambda^2}{m_t^2}}\left(\Delta_{12} - \frac{15}{4} m_t m_b\right)}. \end{aligned} \tag{59}$$

2. and for the P-parity breaking phase:

$$m_1 = 0 \tag{60}$$

$$m_2 \approx 2m_t \tag{61}$$

$$m_3 \approx 2m_b \tag{62}$$

$$m_4 \approx |\sin\delta_0| \sqrt{\frac{15(m_t^2 + m_b^2)}{8\ln\frac{\Lambda^2}{m_t^2}}}. \tag{63}$$

We stress that when $\delta_0 \neq 0$ we have neither scalars no pseudoscalars any longer because the particles which are eigenstates of the energy operator, are mixed of both P-even parity and P-odd parity fields, hence the former classification by parity does not hold for this particular case.

Thus in the Model II the Quasilocal Yukawa interaction with Higgs doublets reduces at low energies to a conventional local one where each Higgs doublet couples to a definite charge current and its *v.e.v.* brings the mass either to *up*- or to *down*- components of fermion doublets. Based on the FCNC suppression, the Model II leads to the relation $m_t >> m_b$ and so to an enhanced coupling of the light scalar (pseudoscalar) boson to the *down*-type quarks while suppressing the coupling to the *up*-type quarks. The Model II has a broad spectrum of excited bound states which is to parametrize the data, in particular, obtained from the Next Linear Collider.

5. Summary

In our talk we have proposed a set of Quasilocal NJL-type quark models (QNJLM) which lead to a larger spectrum of ground and excited states in the polycritical regime. From the viewpoint of the SM, these models are considered as more natural than common extensions of the SM, since they do not enlarge the number of elementary particles in fermionic sector and preserve the symmetries of the SM. In the framework of the QNJLM we have presented two Models which provide at low energies two composite Higgs doublets, as minimal extensions of the Top-Mode Standard Model [11],[14]. In the 2HQ Model I Higgs bosons are rather radial, ground and excited states in the scalar -pseudoscalar channels. In the 2HQ Model II, which consistent with the requirement of natural flavour conservation [16], strong forces lead to the formation of top and bottom bound states (and corresponding condensates) and generate masses of t,b-quarks. In Model II we have concentrated on the scenario where each of the neutral components of the two doublets $\phi_{1,2}$ (with v.e.v. $v_{1,2}$) couple at low energies respectively to the $I_3 = \pm\frac{1}{2}$ fermion fields The FCNC suppression leads to the relation $m_t >> m_b$ and to an enhanced coupling of the light scalar (pseudoscalar) boson to the *down*-type quarks and the charged leptons while suppressing the coupling to the *up*-type quarks. The existence of light neutral Higgs (pseudo)scalar bosons in the framework of 2HDM is not excluded by existing data ($< 40\,GeV$). The chance that it can be seen at the Next Linear Collider in the $\gamma\gamma$ processes has been pointed out in [27],[28]. As a result of complexity of two v.e.v.'s for two composite Higgs doublets the dynamical CP-violation may appear in the Higgs sector. At high energies these channels are strongly coupled and one could say that two-composite Higgs doublets partially represent the mixture with excited states. If such excited states exist then they will modify the Higgs mass predictions. In addition, we remark that low values for the Higgs masses of the additional states could actually change the window for M_H since these states could give a significant contribution to the ρ-parameter [29]. From our consideration we have seen that the appearance of dynamical CP-violation in the Higgs sector imposes strong bounds on Higgs masses, in particular, one light scalar Higgs boson is unavoidable. The experimental implications of such effects are expected to be rather small in the fermion sector of the SM [1],[27]. These effects are observable in decays of heavy Higgs particles (namely, pseudoscalar Higgses may decay into scalar ones, scalar Higgs may decay into pseudoscalar ones) and in decays of Higgses particles into two vector bosons where CP-even and CP-odd amplitudes appear. At high energies the appearance of the appreciable CP-violation could be important both as a source of electron and neutron electric dipole moments [30] and as a mechanism for EW scale baryogenesis[31], [32]. Besides one expect also that modifications of the SM Lagrangian (the Higgs and Top interactions) by higher dimensional vertices may enhance the Higgs production at hadron colliders [33].

6. Acknowledgements

This work has been supported by RFBR (Grant No. 95-02-05346a, No. 98-02-18137), by INTAS (Grant No. 93-283ext), GRACENAS (Grant No. 95-6.3-13) and DFG under contract No. 436 RUS 113/227/1,4.

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