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7 Do Present Experiments Exclude the Fourth Family Quarks as Well as the Existence of More Than One Scalar?

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Abstract. The *spin-charge-family* theory [1–17] predicts the existence of the fourth family to the lower three. It also predicts several scalar fields (the mass eigenstates of the three singlets with the family members quantum numbers and the two triplets with the family quantum numbers) with the weak and the hyper charge of the *standard model* higgs field ($\pm\frac{1}{2}, \mp\frac{1}{2}$, respectively). There is so far no experimental evidence for either the existence of the fourth family quarks with masses below 1 TeV or for the existence of more than one scalar field (however, the Yukawa couplings themselves are the signal that several scalars must exist). If the fourth family quarks have masses above 1 TeV then the experimental evidences [18,19] require that they contribute negligible to either the quark-gluon fusion production of the observed scalar higgs or to the decay of this scalar. It is discussed in this contribution why it is too early to say that the present experiments exclude the fourth family quarks predicted by the *spin-charge-family* theory.

Povzetek. Teorija *spinov-nabojev-družin* [1–17] napove obstoj četrte družine k izmerjenim trem. Napove tudi več skalarnih polj (masnih lastnih stanj treh singletov s kvantnimi števili članov družin in dveh tripletov z družinskimi kvantnimi števili), ki imajo šibki in hiper naboj enak ustreznim nabojem higgsovega polja *standardnega modela* ($\pm\frac{1}{2}, \mp\frac{1}{2}$). Zdi se, da dosedanji poskusi izključujejo obstoj kvarkov četrte družine, z maso pod 1 TeV, pa tudi novih skalarnih polj s tako maso. Če pa naj imajo kvarki četrte družine maso nad 1 TeV, mora biti njihov prispevek k nastanku izmerjenega (higsovega) skalarja, kakor tudi k razpadu tega skalarja, zanemarljiv [18,19]. Prispevek pojasnjuje, zakaj je prezgodaj reči, da četrte družine, ki jo napove teorija *spinov-nabojev-družin* ter novih skalarnih polj (njihov obstoj zagotavljajo že Yukawine sklopitve), ni.

7.1 Introduction

The *spin-charge-family* theory [1–17] predicts before the electroweak break four - rather than the observed three - coupled massless families of quarks and leptons. The 4×4 mass matrices of all the family members demonstrate in this theory the same symmetry [14,15], determined by the scalar fields: the two $\widetilde{SU}(2)$ triplets - the gauge fields of the two family groups operating among families - and the three singlets - the gauge fields of the three charges, $(Q, Q'$ and $Y')$, distinguishing among family members [2,1]. All these scalar fields carry the weak and the hyper charge as does the scalar of the *standard model*: $\pm\frac{1}{2}$ and $\mp\frac{1}{2}$, respectively [17].

Since there is no direct observations of the fourth family quarks masses below 1 TeV, while the fourth family quarks with masses above 1 TeV would contribute, according to the *standard model* (the *standard model* Yukawa couplings of the quarks to the scalar higgs is proportional to $\frac{m_4^\alpha}{v}$, where m_4^α is the fourth family member $(\alpha = u, d)$ mass and v the vacuum expectation value of the scalar), to either the quark-gluon fusion production of the scalar field (the higgs) or to the scalar field decay, too much in comparison with the observations, the high energy physicists do not expect the existence of the fourth family members at all [18,19].

Does this mean that there does not exist the fourth family coupled to the observed three?

Before discussing the question to which extent can be the theoretical interpretations of the experimental data, grounded on the *standard model* assumptions, acceptable for four families, while they are obviously working well for three families, let be pointed out what supports the *spin-charge-family* theory to be the right next step beyond the *standard model*. This theory is only able not to explain - while starting from the very simple action in $d \geq (13+1)$, Eqs. (7.8, 7.9) of Sect. 7.4, with massless fermions with the spin of the two kinds (one kind taking care of the spin and of the charges of the family members the second kind taking care of the families (Eq. (6.50))), which couple only to the gravity (through the vielbeins and the two kinds of the corresponding spin connections (Eqs. (7.8, 7.9))) - all the assumptions of the *standard model*, but also to answer several open questions beyond the *standard model*. It offers the explanation for [1–17]:

a. the appearance of all the charges of the left and right handed family members and for their families and their properties,

b. the appearance of all the corresponding vector and scalar gauge fields and their properties (explaining the appearance of the higgs and the Yukawa couplings),

c. the appearance and properties of the dark matter,

d. the appearance of the matter/antimatter asymmetry in the universe.

The theory predicts for the low energy regime:

i. The existence of the fourth family to the observed three.

ii. The existence of twice two triplets and three singlets of scalars, all with the properties of the higgs with respect to the weak and hyper charges, explaining the existence of the Yukawa couplings. Besides the higgs also a few of the others will be observed at the LHC.

iii. There are several other predictions.

Since the experimental accuracy of the $(3 \times 3$ submatrix of the $4 \times 4)$ mixing matrices is not high enough, it is not yet possible to estimate masses of the fourth family members by fitting the experimental data to the parameters of mass matrices, determined by the symmetry as predicted by the *spin-charge-family* [15,14]. While the fitting procedure is not influenced considerably by the accuracy of the measured masses of the lower three families, the accuracy of the measured values of mixing matrices do influence, as expected the fitting results very much. The fact that the fourth family quarks have not yet been observed - directly or indirectly - pushes the fourth family quarks masses to ≈ 2 TeV or higher.

The more effort and work is put into the *spin-charge-family* theory, the more explanations of the observed phenomena and the more predictions for the future observations follow out of it. Offering the explanation for so many observed phenomena - keeping in mind that all the explanations for the observed phenomena originate in a simple starting action - qualifies the *spin-charge-family* theory as the candidate for the next step beyond the *standard model*.

Since in the *spin-charge-family* theory all the low energy degrees of freedom of elementary fields - fermions and vector and scalar bosons - follow from a simple starting action, and since also the dynamics is determined in the starting action, it would in principle be possible to calculate all the properties of the fields in the low energy regime, if we would know the boundary conditions. This is, of course, too ambitious program, not only because the boundary conditions are not known, but also because of too many degrees of freedom of fermions and bosons, in particular at phase transitions (as teaches us physics of fluids and condensed matter).

This paper argues for the existence of the fourth family and of several scalar fields predicted by the *spin-charge-family* theory, discussing the arguments why the contribution of the fourth family quarks to the quark-gluon fusion at LHC, as well as the contribution of this family to the decay of the Higgs's scalar might not disagree with the observations, as long as the interpretations of the events rely on the *standard model* assumptions Ref. [20], which are not in agreement with the *spin-charge-family* theory.

Sect. 7.2 discusses the arguments why the fourth family might exist although has not yet been observed - directly or indirectly.

The *spin-charge-family* theory is presented in the main talk of the author of this contribution.

7.2 The fourth family in the *spin-charge-family* theory and the experimental constraints against it

The *spin-charge-family* theory predicts the fourth family to the observed three. The calculation of the fourth family properties to the observed three, when taking into account the symmetry of mass matrices predicted by this theory and fitting the consequently allowed parameters of mass matrices to the experimental data, shows [14,15], that the measured matrix elements of the 3×3 - submatrices of the 4×4 - mixing matrices are far from being accurate enough even for quarks to determine masses of the fourth family members. In Subsect. 7.2.3 a short report on this calculation is presented. More can be found on Refs. [27,1,15] and the references cited there.

Since there has been no direct observation of the fourth family quarks with the masses below 1 TeV, while the *standard model* without the fourth family is in much better agreement with the experiments than with the fourth family included, the high energy physicists do not expect the existence of the fourth family members at all [18,19].

Should the explanation of the so far obtained experimental data by using the *standard model* assumptions be accepted as the definite experimental evidence that there are only three families and that the fourth family of quarks and leptons does not exist?

Let us try to understand how far are the interpretations of the experimental data trustworthy, if following more or less the *standard model* assumptions:
i. The *standard model* assumes one scalar doublet, the higgs (there are also models using two scalar doublets or more, all more or less following the idea of the *standard model* higgs [21,22]).
ii. In calculations the validity of the Yukawa couplings in the perturbation calculations are assumed.
iii. Calculations have been done in the leading order, next to the leading order and even some in one order more.
iv. The 4×4 mixing matrices elements of the fourth family members to the observed three were just assumed (or neglected).

The review article [20] discusses, reporting on many papers (≈ 200), a possibility of the existence of the fourth family due to the experimental data and theoretical analyses of the data, presenting also assumptions on which the theoretical analyses were done. The author discusses the (non)existence of the fourth family quarks and leptons due to direct searches for the fourth family members, due to changes in mixing matrices if there exists the fourth family, due to measurements testing the electroweak precision data with and without the existence of the fourth family, and in particular due to the analyzes of the higgs boson production and decay. The author concludes pointing out that if taking seriously that there exists only one scalar doublet and if assuming perturbativity of the Yukawa couplings and the Dirac mass of the heavy neutrino, "then the fourth family of fermions can not accommodate the data for the higgs searches" and its decay.

This review article [20] appeared in 2013. The new data [24], reported also in the review talk [25], are not in contradiction with the conclusions of Ref. [20], while the measured mixing matrix elements for quarks - averaged over data of several experimental groups - are still far from being accurate enough to allow the *spin-charge-family* theory to predict the fourth family quarks masses.

Although the assumptions, used to analyze the experimental data, might seem to most of high energy physicists acceptable, the assumptions do not appear so trustworthy when looking at them from the point of view of the *spin-charge-family* theory, what it will be done in this section.

Let us point out the differences between the generally used assumptions in the analyses of the experimental data searching for new scalars and new family and the properties of the fermions and scalar fields in the *spin-charge-family* theory.

Most commonly accepted assumptions [20,24]:
A. There is only one Higgs doublet. If there are more, their properties (their Lagrange function) resemble the properties of the *standard model* higgs.
B. The Yukawa couplings are used in the higher order corrections.
C. The Dirac neutrino masses is used.
D. The perturbativity of the theory is assumed.

Arguments against these common assumptions:
A′. Already the higgs enters into the *standard model* by "hand", with assumed Lagrange function which couples the higgs to the weak bosons ($W^{\pm}_{m}$ and Z^{0}_{m}) and "dresses" fermions with the (appropriate) weak and hyper charges. The higgs does not carry the family quantum numbers. To our understanding, just repeating the "game of the higgs" for several higges without a deeper understanding of the origin of scalars can hardly be the right way beyond the *standard model*.
B′. The Yukawa couplings are also put by "hand" into the *standard model* to compensate the higgs independence of the family quantum numbers. Using the Yukawa couplings in perturbative way might have no theoretical support in calculations with corrections next to leading and next to next to leading orders.
C′. The Dirac neutrino masses seems natural, since all the other family members do have the Dirac masses, but either the Dirac or the Majorana mass of neutrinos must be grounded in a deeper understanding of the origin of fermion masses.
D′. The perturbativity of the theory, originating in effective Lagrange function, at and around the phase transition of the electroweak break, is also questionable, since the acceptance of the effective theory might easily break down.

Calculations done under the assumptions presented from A.-D. lead, due to Ref. [20], to the conclusion:
i. The ratio of the gluon-gluon fusion generating the higgs and decaying into two ZZ if taking into account the four families or only three is $\approx 5-8$.
ii. The ratio of the gluon-gluon fusion generating higgs and decaying into $b\bar{b}$ if taking into account the four families or only three is ≈ 5.
iii.The most stringent is, due to Ref. Lenz, the predicted underproduction (for a factor of 5) in the two γ's channel, if the fourth family is included into calculations with respect to the calculations with only three families.
The author of Ref. [20] reports also some additional drawbacks (calling them minor) of the theoretical interpretation of experiments, like: not taking into account the change of the mixing matrix if there are four families instead of three, not allowing a large enough interval the masses of the fourth family, not taking into account the decays of higgs through the fourth family neutrinos if they appear to be light enough (smaller than $\frac{m_H}{2}$).

The spin-charge-family theory is disagrees with the assumptions A.- D.:
A″. In the *spin-charge-family* theory there are three singlet and two triplet scalar fields originating like all the gauge fields (vectors, tenzors and scalars) in the starting action (this can be read in Eqs. (7.8, 7.9) and in the talk of N.S.M.B. in this Proc., Eqs. (19,20)) as the gauge spin connection fields and manifesting in $(3+1)$ as scalars with the space index $s = (7, 8)$, all carrying the weak and the hyper charges (determined by the space index $s = (7, 8)$) of the *standard model* higgs (in the talk of N.S.M.B. in this Proc., Eq. (21)). The three singlets carry besides the higgs quantum numbers also the family members quantum numbers (Q, Q', Y'), the two triplets carry besides the higgs quantum numbers also the family quantum numbers, Eq. (5) in the talk of N.S.M.B. in this Proc..
B″. Scalars start as massless gauge fields, gaining masses when interacting with the condensate (Table 1. in the talk of N.S.M.B. in this Proc.) and change their properties when obtaining at the electroweak break the nonzero vacuum expectation

values. Each family member of each family couples to a different superposition of the scalar mass eigenstates, what correspondingly determines the Yukawa couplings (Eqs. (25-28) in the talk of N.S.M.B. in this Proc.). Since it is not known at which scale does the electroweak break occur, the perturbativity of the Yukawa couplings might not be an acceptable assumption, in particular since at phase transitions all the systems manifest the long range properties.
C″. In the *spin-charge-family* theory all the family members have the Dirac masses. However, to the mass matrices of each family member besides the scalar fields carrying the family quantum numbers - the two triplets - also the scalars carrying the family members quantum numbers - the three singlets (Q, Q', Y') - contribute. To neutrino masses besides the two triplets only the singlet - the gauge field of Y' - contribute. All these contributions are highly nonperturbative [12–15]. The three singlets contribute to the off diagonal mass matrix elements only in the higher order corrections, what makes masses of the family members so different.
D″. Scalars with the space index $s = (7, 8)$ (the two triplets and the three singlets) gain in a highly nonperturbative way nonzero vacuum expectation values, obeying after the electroweak break the approximate effective Lagrange function (Eq. (24)), for which one can't expect that close to the phase transition the perturbativity would work.

Let us look at the properties of the scalar fields contributing to the masses of the lower four families and of the $W^{\pm}_{m}, Z^{0}_{m}$ in the *spin-charge-family* theory in more details.

7.2.1 Scalar fields contributing to the mass matrices of the lower four families in the *spin-charge-family* theory

The *spin-charge-family* theory predicts twice (almost) decoupled four families in the low energy regime (Refs. [1,3,2,26]. We discuss here only properties of the lower four families.

To understand better how do in the *spin-charge-family* theory the scalar fields determine the properties of families of each family member - after the loop corrections are taken into account in all orders - the Lagrange density of the fermion mass term, Eq. (7.9) and Eq. (20) in the talk of N.S.M.B. in this Proc., the quarks part in particular, is rewritten so that (massless) quark states, u^k and d^k, k is the family index, enter explicitly into expressions.

$$
\begin{aligned}
&\mathcal{L}_{f=qm} = \\
&\frac{1}{2}\{[\psi^{\dagger}_{L}\gamma^0(\sum_{A,i,+,-} \overset{78}{(\pm)}\, \tau^{Ai} g^{Ai} A^{Ai}_{\pm})\, \psi_R] + [\psi^{\dagger}_{L}\gamma^0(\sum_{A,i,+,-} \overset{78}{(\pm)}\, \tau^{Ai} g^{Ai} A^{Ai}_{\pm})\, \psi_R]^{\dagger}\} = \\
&= \frac{1}{2}\sum_{k,l}\{[u^{k\dagger}_{L}\gamma^0(\sum_{A,i} \overset{78}{(-)}\, \tau^{Ai} g^{Ai} A^{Ai}_{-})\, u^{l}_{R}] + [u^{k\dagger}_{L}\gamma^0(\sum_{A,i} \overset{78}{(-)}\, \tau^{Ai} g^{Ai} A^{Ai}_{-})\, u^{l}_{R}]^{\dagger} + \\
&\quad [d^{k\dagger}_{L}\gamma^0(\sum_{A,i} \overset{78}{(+)}\, \tau^{Ai} g^{Ai} A^{Ai}_{+})\, d^{l}_{R}] + [d^{k\dagger}_{L}\gamma^0(\sum_{A,i} \overset{78}{(+)}\, \tau^{Ai} g^{Ai} A^{Ai}_{+})\, d^{l}_{R}]^{\dagger}\}. \quad (7.1)
\end{aligned}
$$

Operators $\tau^{Ai} = \sum_{s,t} c^{Ai}{}_{st}\, S^{st}$ are defined in Eqs. (3,4) in the talk of N.S.M.B. in this Proc. and in Eqs. (7.10, 7.11), scalar fields $A^{Ai}_{\pm} = \sum_{s,t} c^{Ai}{}_{st}\, \omega^{st}{}_{\pm}$ and

$\gamma^0 \stackrel{78}{(\pm)}$ are defined in Eqs. (19-21) in the talk of N.S.M.B. in this Proc.. The coupling constants of the two triplets ($\tilde{A}^{\tilde{1}i}_{\pm}$, $\tilde{A}^{i}_{\tilde{N}_L \pm}$) and three singlets ($A^{Q}_{\pm}$, $A^{Q'}_{\pm}$, $A^{Y'}_{\pm}$) are in Eq. (7.1) written explicitly.

Operators τ^{Ai} and $\gamma^0 \stackrel{78}{(\pm)}$ are Hermitian. In what follows it is assumed that the scalar fields A^{Ai}_{s} are Hermitian as well and consequently it follows $(A^{Ai}_{\pm})^{\dagger} = A^{Ai}_{\mp}$). While the family operators $\tilde{\tau}^{1i}$ and $\tilde{N}^{i}_{L}$ commute with $\gamma^0 \stackrel{78}{(\pm)}$, the three family members operators (Q, Q', Y') do not, but one sees that

$$
\begin{aligned}
(u^{k\dagger}_{L} \gamma^0 \sum_{A,i,+,-} (Q, Q', Y')\, g^{(Q,Q',Y')} \stackrel{78}{(-)} A^{(Q,Q',Y')}_{-} u^{l}_{R})^{\dagger} = \\
u^{l\dagger}_{R} \gamma^0 \,(Q, Q', Y')\, g^{(Q,Q'Y')} \stackrel{78}{(+)} A^{(Q,Q',Y')\dagger}_{-} u^{k}_{L} \\
= u^{l\dagger}_{R} \,(Q^{k}_{R}, Q'^{k}_{R}, Y'^{k}_{R})\, g^{Q,Q',Y'} A^{(Q,Q',Y')}_{+} u^{k}_{R}\,, \quad (7.2)
\end{aligned}
$$

where $(Q^{k}_{R}, Q'^{k}_{R}, Y'^{k}_{R})$ denote the eigenvalues of the spinor state u^{k}_{R}.

The off diagonal matrix elements of mass matrices, Eq. (7.7), start to be dependent on the family members quantum numbers only in loop corrections, consequently the contributions of the loop corrections to all orders are indeed important.

Couplings of u_k and d_k to the scalars carrying the family members quantum numbers are determined also by the eigenvalues of the operators (Q, Q', Y') on the family members states. Strong influences of the scalar fields carrying the family members quantum numbers on the masses of the lower (observed) three families of quarks manifest in huge differences in masses of u_k and d_k, $k = (1, 2, 3)$, among family members (u, d). For the fourth family quarks, which are more and more decoupled from the observed three families the higher are their masses [15,14], the influence of the scalar fields carrying the family members quantum numbers on their masses is expected to be much weaker. *Correspondingly might become the* u_4 *and* d_4 *masses closer to each other the higher are their masses and the weaker is their couplings* (the mixing matrix elements) *to the lower three families.*

The superposition of the scalar eigenstates which couple to the fourth family quarks might therefore differ a lot from those which couple to the lower three families.

Although the gluons couple in the gluon-gluon fusion to all the quarks in an equivalent way, yet *the family members with different family quantum number contribute to the production of different scalar mass eigenstates differently*, which might not be in agreement with the simple *standard model* prediction, that the fourth family couples to the observed higgs proportionally to their masses ($\frac{m_i}{v}$).

If the masses of the fourth family quarks are close to each other, then u_4 *and* d_4 *contribute in the quark-gluon fusion very little to the production of the observed scalar field - the higgs - if the higgs is a superposition of different scalar fields mass eigenstates then the scalar, to which the fourth family quarks mostly couple, as it is expected.*

The scalar fields from the starting action to the effective action Let us discuss the scalar fields, which contribute to the electroweak break with nonzero vacuum expectation values, from the point of view of the starting action (Eq. (7.8)) in order to try to understand better their properties at the electroweak break.

The action (Eq. (7.8)) manifests that there are only $\tilde{\omega}_{mn\sigma}$, which are coupled to the vector gauge fields $\mathcal{A}^{Ai}_{m}$ ([27], Eqs. (19,20) in [28], as well as Subsect. 6.3.2, page 88, in the talk of N.S.M.B. in this Proc.) on the tree level.

The vector gauge fields $\mathcal{A}^{Ai}_{m}$, namely, appear in the action, Eq. (7.8), as (Subsect. 3.2. in my talk in this Proc.)

$$f^{\sigma}{}_{m} = \sum_{A} \vec{\tau}^{A\sigma}\, \vec{\mathcal{A}}^{A}_{m}\,, \tag{7.3}$$

where $\tau^{Ai} = \sum_{st} c^{Ai}{}_{st}\, M^{st}$, $M^{st} = S^{st} + L^{st}$, $\{\tau^{Ai}, \tau^{Bj}\}_{-} = i f^{Aijk} \tau^{Ak}\, \delta^{AB}$, $\vec{\tau}^{A} = \vec{\tau}^{A\sigma}\, p_{\sigma} = \vec{\tau}^{A\sigma}{}_{\tau}\, x^{\tau}\, p_{\sigma}$, $\tau^{Ai\sigma} = \sum_{st} c^{Ai}{}_{st}\, M^{st\sigma} = \sum_{st} c^{Ai}{}_{st}\, (e_{s\tau}\, f^{\sigma}{}_{t} - e_{t\tau}\, f^{\sigma}{}_{s}) x^{\tau}$, $\mathcal{A}^{Ai}_{m} = \sum_{st} c^{Ai}{}_{st}\, \omega^{st}{}_{m}$. The relation between $\omega^{st}{}_{m}$ and vielbeins is determined by Eq. (7.4), if there are no spinor sources present.

Correspondingly the vector gauge fields gain masses on the tree level through $f^{\sigma}{}_{m} f^{\tau}{}_{n}\, (\tilde{\omega}^{mn}{}_{[\sigma,\tau]} - \tilde{\omega}^{m}{}_{m'[\tau} \tilde{\omega}^{m'n}{}_{\sigma]})$ ([ ] means that the two indices must be exchanged), when at the electroweak break the two triplets and the three singlets gain the nonzero vacuum expectation values. Indeed only one superposition - one triplet - is involved. Since only $f^{\sigma}{}_{m}$ represents the vector gauge fields ($f^{\mu}{}_{s} = 0$ [1,28]) in the low energy regime, all the rest of scalar fields (the second triplet and the three singlets) contribute to masses of $W^{\pm}_{m}$, Z^{0}_{m} only in loop corrections.

The action (Eq. (7.8)) leads to the equations of motion (Ref. [1], Eqs. (31,32))

$$\begin{aligned}
0 &= 2\alpha \left[f^{\beta}{}_{b}\, R^{ba}{}_{[\beta\alpha]} - \frac{1}{2}\, e^{a}{}_{\alpha}\, R \right] \\
&\quad + 2\tilde{\alpha} \left[f^{\beta}{}_{b}\, \tilde{R}^{ba}{}_{[\beta\alpha]} - \frac{1}{2}\, e^{a}{}_{\alpha}\, \tilde{R} \right] \\
&\quad + \bar{\Psi} \gamma^{a} p_{0\alpha} \Psi - f^{\beta}{}_{b} e^{a}{}_{\alpha} \left(p_{\beta} (\bar{\Psi} \gamma^{b} \Psi) - p_{\alpha} (\bar{\Psi} \gamma^{a} \Psi) \right)\,, \\
p_{0\alpha} &= p_{\alpha} - \frac{1}{2}\, S^{cd}\, \omega_{cd\alpha} - \frac{1}{2}\, \tilde{S}^{cd}\, \tilde{\omega}_{cd\alpha}\,, \\
R^{ab}{}_{[\alpha\beta]} &= \partial_{[\alpha}\, \omega^{ab}{}_{\beta]} + \omega^{a}{}_{c[\alpha} \omega^{cb}{}_{\beta]}\,, \\
\tilde{R}^{ab}{}_{[\alpha\beta]} &= \partial_{[\alpha}\, \tilde{\omega}^{ab}{}_{\beta]} + \tilde{\omega}^{a}{}_{c[\alpha} \tilde{\omega}^{cb}{}_{\beta]}\,,
\end{aligned} \tag{7.4}$$

$$\begin{aligned}
f^{\alpha}{}_{c}\, \omega_{[a}{}^{c}{}_{b]} + f^{\alpha}{}_{[a}\, \omega_{b]}{}^{c}{}_{c} &= \frac{1}{E}\, \partial_{\beta} (E f^{\alpha}{}_{[a} f^{\beta}{}_{b]}) + \frac{1}{2}\, \bar{\Psi}\, f^{\alpha}{}_{c} \gamma^{c}\, S_{ab} \psi\,, \\
f^{\alpha}{}_{c}\, \tilde{\omega}_{[a}{}^{c}{}_{b]} + f^{\alpha}{}_{[a}\, \tilde{\omega}_{b]}{}^{c}{}_{c} &= \frac{1}{E}\, \partial_{\beta} (E f^{\alpha}{}_{[a} f^{\beta}{}_{b]}) + \frac{1}{2}\, \bar{\Psi}\, f^{\alpha}{}_{c} \gamma^{c}\, \tilde{S}_{ab} \psi\,.
\end{aligned} \tag{7.5}$$

One can read in Eqs. (7.4, 7.5) the interactions among the gauge fields and the interactions of the gauge fields with the fermion fields (in particular we point out the condensate [1] and Table 6.1 (on page 83) in the N.S.M.B. talk in this Proc.).

The appearing of the condensate, its interaction with the scalars and the behavior of scalars at the electroweak phase transition are expected to be highly nonperturbative effects. It is assumed so far (estimating very roughly the degrees of freedom and the interactions among scalars and among scalars and fermions) that the effective Lagrange density of scalars, contributing to the electroweak break, might after the phase transition manifest the *standard model* assumptions, changing from the starting action Lagrange density Eq. (7.8) $\mathcal{L}_s = E\{(p_m A^{Ai}_s)^\dagger (p^m A^{Ai}_s) - (m'_{Ai})^2 A^{Ai\dagger}_s A^{Ai}_s\}$ to

$$\mathcal{L}_{sg} = E \sum_{A,i} \{(p_m A^{Ai}_s)^\dagger (p^m A^{Ai}_s) - (-\lambda^{Ai} + (m'_{Ai})^2))\, A^{Ai\dagger}_s A^{Ai}_s + \sum_{B,j} \Lambda^{AiBj} A^{Ai\dagger}_s A^{Ai}_s A^{Bj\dagger}_s A^{Bj}_s\}, \tag{7.6}$$

where $-\lambda^{Ai} + m'^2_{Ai} = m^2_{Ai}$ and m_{Ai} manifests as the mass of the A^{Ai}_s scalar.

Whether or not this is an acceptable effective Lagrange function or not remains to be proved.

7.2.2 The contribution of the fourth family quarks of equal masses to the production of the scalar fields

In the N.S.M.B. talk in this proceedings, Subsect. 6.4.2 (page 101), a possibility is discussed that if the fourth family quarks have approximately equal masses (the u_4-quarks and d_4-quarks might have similar masses, if their masses are mostly determined by the scalars with the family quantum numbers, as discussed in Subsect. 7.2.1), while u_4 and d_4 couple to the scalar fields determining their masses with the opposite phases, the contribution of the fourth family quarks to the production of scalar fields in the gluon-gluon fusion can be negligible.

Since the family quantum numbers ($\tilde{\tau}^1$, $\tilde{N}^L$) commute with the weak and the hyper charges, the scalar fields carrying the family quantum numbers, $\vec{\tilde{A}}^{1,N_L}_{\substack{78\\(\pm)}}$, distinguish among u_4 and d_4 only due to the operator $\overset{78}{(\pm)}$. Couplings of u_4 and d_4 to those scalar fields, which carry in addition to the weak and the hyper charge the family members quantum numbers - to the three singlets ($A^{Q}_{\substack{78\\(\pm)}}, A^{Q'}_{\substack{78\\(\pm)}}, A^{Y'}_{\substack{78\\(\pm)}}$) - depend on the eigenvalues of (Q, Q', Y') on the quark states, which are different for u_i and d_i quarks.

Since the masses of u_4 and d_4 are only approximately equal, the fourth family quarks can still weakly contribute to the production of the scalar fields, in particular to those which mostly determine masses of the fourth family members.

7.2.3 Mass matrices of family members and the masses of the fourth family quarks

The *spin-charge-family* theory [2,1,14,15] predicts the mass matrices of the family members α for each groups of four families, Eq. (7.7).

$$\mathcal{M}^\alpha = \begin{pmatrix} -a_1 - a & e & d & b \\ e & -a_2 - a & b & d \\ d & b & a_2 - a & e \\ b & d & e & a_1 - a \end{pmatrix}^\alpha . \tag{7.7}$$

The mass matrices are determined at the electroweak break, when the scalar fields with the space index $s = (7, 8)$ (the three singlets carrying the family members quantum numbers and the two triplets carrying the family quantum numbers, the two triplets and the three singlets interacting among themselves, Eq. (7.6)) get nonzero vacuum expectation values. In loop corrections the singlets influence all the matrix elements of each mass matrix while keeping the $\widetilde{SU}_{\widetilde{SO}(3,1)}(2) \times \widetilde{SU}_{\widetilde{SO}(4)}(2) \times U(1)$ symmetry unchanged.

In Refs. [14,15] the twice 6 parameters of the two mass matrices of the lower group of four families of u and d quarks, in general non Hermitean, presented in Eq. (7.7), were fitted to:
i. twice three masses of the u_i, $i = 1, 2, 3$ (u, c, t) and d_i, $i = 1, 2, 3$ (d, s, b) quarks and
ii. 3×3 submatrix of the 4×4 quark mixing matrix. Although the $(n-1) \times (n-1)$ submatrix of the $n \times n$ unitary matrix, if accurately known, determines uniquely the $n \times n$ matrix for $n \geq 4$, we were not able to determine the masses of the fourth families, even not after assuming that the mass matrices are real, since the 3×3 submatrices are not known accurately enough. We only could tell the fourth family matrix elements of the mixing matrix after assuming the masses of the fourth family quarks.

It turned out that for the masses of the fourth family quarks above 1 TeV the mass matrices are more and more democratic and the fourth family quarks are more and more decoupled from the lower three families the larger are the fourth family masses. It correspondingly appears that the masses of u_4 and d_4 are closer to each other the smaller is contribution of the scalar fields with the family members quantum numbers to their masses. The results are presented in Ref. [15].

7.3 Concluding remarks

In this contribution the arguments against the conclusions of most high energy physicists that present experiments can hardly leave any room for the existence of the fourth family members is discussed.

The analysis of experiments, which are based on the assumptions of the *standard model* - i. on the existence of one scalar doublet, if there are several they follow properties of the higgs, ii. on the perturbativity of the theory, iii. on guessing the mixing matrices elements of the fourth family members to the observed three - might from the point of view of the *spin-charge-family* theory not be acceptable.

The main arguments against the *standard model* assumptions through the "eyes" of the *spin-charge-family* are:
i. Assuming the existence of one scalar fields, or even several scalars repeating the idea of the *standard model* higgs, is too restrictive. In the *spin-charge-family* theory there are three singlet and two triplet scalar fields, which all originate (like all the gauge fields and the gravity in $(3+1)$ do) in the starting action as the gauge spin connection fields and manifesting in $(3+1)$ as scalars with the space index $s = (7, 8)$, all carrying the weak and the hyper charges (determined by the space index $s = (7, 8)$) of the *standard model* higgs. The three singlets carry besides the higgs quantum numbers also the family members quantum numbers (Q, Q', Y'), the two triplets carry besides the higgs quantum numbers also the family quantum

numbers.
ii. Scalars start as massless gauge fields, gaining masses when interacting with the condensate and changing their properties when at the electroweak break when gaining the nonzero vacuum expectation values. Each family member of each family couples to a different superposition of the scalar mass eigenstates, what correspondingly determines the Yukawa couplings. All these contributions are highly nonperturbative.
In addition, the three singlets contribute to the off diagonal mass matrix elements only in the higher order corrections, what makes masses of the family members so different. Since it is also not known at which scale does the electroweak break occur, the perturbativity might not be an acceptable assumption even if at the low enough energies the effective Lagrange density behaves perturbatively.
iii. Also the mixing matrix elements from the fourth family members to the the rest three might influence considerably the interpretation of the experimental data. Also since each family member couples to different superposition of the scalar mass eigenstates.
Let us add that each family member can couple to the scalar fields with its own phase. In the case that the fourth family quark u_4 couples to the scalar fields with the opposite phase than the d_4 quark, and that their masses are closed to each other, what seems to be the case in the *spin-charge-family* theory, then the fourth family quarks contribution to the production of higgs and its decay might be very small, also since the superposition of the scalar mass eigenstates which couple to the fourth family quarks differ a lot from those which couple to the lower three families.

Although the gluons couple in the gluon-gluon fusion to all the quarks in an equivalent way, yet the family members with different family quantum number contribute to the production of different scalar mass eigenstates differently, what means that the simple *standard model* prediction, that the fourth family couples to the observed higgs proportionally to their masses ($\frac{m_i}{v}$), is not acceptable.

Let us point out at the end that the ability of the *spin-charge-family* theory, which starts with a simple action with fermions carrying two kinds of spins and no charges in $d > (3+1)$ and interacting with only gravitational field, to offer the explanation **a.i.** for all the assumptions of the standard model, **a.ii.** for the appearance of the family members and the families, **a.iii.** for the appearance of the gauge vector fields and their properties, **a.iv.** for the appearance of the scalar fields explaining the higgs and the Yukawa couplings, **a.v.** for the appearance of the dark matter, **a.vi.** for the appearance of the matter/antimatter asymmetry in the universe, suggests that this theory must be taken as a candidate showing next step beyons the *standard model*. Correspondingly must the prediction of this theory that there exists the fourth family coupled to the observed three and that there exist several scalar fields, which explain besides the origin of the higgs also the Yukawa couplings, be taken seriously.

7.4 Appendix: Spin-charge-family theory, action and assumptions

I present in this appendix the *assumptions* of the *spin-charge-family* theory, on which the theory is built - following a lot the equivalent sections in Refs. [2,1] - starting with the simple action for fermions and the gravity fields.

A i. In the action [2–4,1], Eq. (7.8), fermions ψ carry in $d = (13+1)$ as the *internal degrees of freedom only two kinds of spins* (no charges), which are determined by the two kinds of the Clifford algebra objects (there exist no additional Clifford algebra objects (6.48)) - γ^a and $\tilde{\gamma}^a$ - and *interact correspondingly with the two kinds of the spin connection fields* - $\omega_{ab\alpha}$ and $\tilde{\omega}_{ab\alpha}$, the *gauge fields* of $S^{ab} = \frac{i}{4}\,(\gamma^a\gamma^b - \gamma^b\gamma^a)$ (the generators of $SO(13,1)$) and $\tilde{S}^{ab} = \frac{i}{4}\,(\tilde{\gamma}^a\tilde{\gamma}^b - \tilde{\gamma}^b\tilde{\gamma}^a)$ (the generators of $\widetilde{SO}(13,1)$) - and the *vielbeins* $f^{\alpha}{}_{a}$.

$$
\begin{aligned}
\mathcal{A} &= \int d^dx\, E\, \mathcal{L}_f + \int d^dx\, E\, (\alpha R + \tilde{\alpha}\tilde{R}),\\
\mathcal{L}_f &= \frac{1}{2}\,(\bar{\psi}\,\gamma^a p_{0a}\psi) + h.c.,\\
p_{0a} &= f^{\alpha}{}_a p_{0\alpha} + \frac{1}{2E}\{p_\alpha, Ef^{\alpha}{}_a\}_-, \quad p_{0\alpha} = p_\alpha - \frac{1}{2}S^{ab}\omega_{ab\alpha} - \frac{1}{2}\tilde{S}^{ab}\tilde{\omega}_{ab\alpha},\\
R &= \frac{1}{2}\{f^{\alpha[a}f^{\beta b]}\,(\omega_{ab\alpha,\beta} - \omega_{ca\alpha}\,\omega^{c}{}_{b\beta})\} + h.c.\,,\\
\tilde{R} &= \frac{1}{2}\{f^{\alpha[a}f^{\beta b]}\,(\tilde{\omega}_{ab\alpha,\beta} - \tilde{\omega}_{ca\alpha}\,\tilde{\omega}^{c}{}_{b\beta})\} + h.c.\,.
\end{aligned}
\tag{7.8}
$$

Here [1] $f^{\alpha[a}f^{\beta b]} = f^{\alpha a}f^{\beta b} - f^{\alpha b}f^{\beta a}$. R and $\tilde{R}$ are the two scalars (the two curvatures) [2].

A ii. The manifold $M^{(13+1)}$ breaks first into $M^{(7+1)}$ times $M^{(6)}$ (manifesting as $SO(7,1)\times SU(3)\times U(1)$), affecting both internal degrees of freedom - the one represented by (the superposition of) S^{ab} and the one represented by (the superposition of) $\tilde{S}^{ab}$. Since the left handed (with respect to $M^{(7+1)}$) spinors couple differently to scalar (with respect to $M^{(7+1)}$) fields than the right handed ones, the break can leave massless and mass protected $2^{((7+1)/2-1)}$ massless families (which decouple into twice four families). The rest of families get heavy masses [3].

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[1] $f^{\alpha}{}_a$ are inverted vielbeins to $e^a{}_\alpha$ with the properties $e^a{}_\alpha f^{\alpha}{}_b = \delta^a{}_b$, $e^a{}_\alpha f^{\beta}{}_a = \delta^{\beta}_{\alpha}$, $E = \det(e^a{}_\alpha)$. Latin indices $a, b, .., m, n, .., s, t, ..$ denote a tangent space (a flat index), while Greek indices $\alpha, \beta, .., \mu, \nu, ..\sigma, \tau, ..$ denote an Einstein index (a curved index). Letters from the beginning of both the alphabets indicate a general index ($a, b, c, ..$ and $\alpha, \beta, \gamma, ..$), from the middle of both the alphabets the observed dimensions $0, 1, 2, 3$ ($m, n, ..$ and $\mu, \nu, ..$), indices from the bottom of the alphabets indicate the compactified dimensions ($s, t, ..$ and $\sigma, \tau, ..$). We assume the signature $\eta^{ab} = diag\{1, -1, -1, \cdots, -1\}$.

[2] R and $\tilde{R}$ are expressible with vielbeins and their derivatives, when there are no fermions present [1,34].

[3] A toy model [29,30] was studied in $d = (5+1)$ with the same action as in Eq. (7.8). The break from $d = (5+1)$ to $d = (3+1)\times$ an almost S^2 was studied. For a particular choice of vielbeins and for a class of spin connection fields the manifold $M^{(5+1)}$ breaks

A iii. The manifold $M^{(7+1)}$ breaks further into $M^{(3+1)}\times M^{(4)}$.

A iv. The scalar condensate (Table 7.1) of two right handed neutrinos with the family quantum numbers of one of the two groups of four families, brings masses of the scale of the unification ($\approx 10^{16}$ GeV or higher) to all the vector and scalar gauge fields, which interact with the condensate [2].

A v. There are nonzero vacuum expectation values of the scalar fields with the space index $s = (7,8)$, conserving the electromagnetic and colour charge, which cause the electroweak break and bring masses to all the fermions and to the heavy bosons.

Comments on the assumptions:

C i. The starting action contains all degrees of freedom, either for fermions or for bosons, needed to manifest at low energy regime in $d = (3+1)$ all the vector and scalar gauge fields and the one family members as well as families of quarks and leptons as assumed by the *standard model*: **a.** One representation of $SO(13,1)$ contains, if analyzed with respect to the *standard model* groups ($SO(3,1)\times SU(2)\times U(1)\times SU(3)$) all the members of one family (Table 6.4, page 89), left and right handed, quarks and leptons (the right handed neutrino is one of the family members), anti-quarks and anti-leptons, with the quantum numbers required by the *standard model* [4]. **b.** The action explains the appearance of families due to the two kinds of the infinitesimal generators of groups: S^{ab} and $\tilde{S}^{ab}$ [5]. **c.** The action explains the appearance of the gauge fields of the *standard model* [2,1] [6]. **d.** It explains the appearance of the scalar higgs and Yukawa couplings [7].

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into $M^{(3+1)}$ times an almost S^2, while $2^{((3+1)/2-1)}$ families remain massless and mass protected. Equivalent assumption, its proof is in progress, is made in the $d = (13+1)$ case.

4 It contains the left handed weak ($SU(2)_I$) charged and $SU(2)_{II}$ chargeless colour triplet quarks and colourless leptons (neutrinos and electrons), and the right handed weak chargeless and $SU(2)_{II}$ charged coloured quarks and colourless leptons, as well as the right handed weak charged and $SU(2)_{II}$ chargeless colour anti-triplet anti-quarks and (anti)colourless anti-leptons, and the left handed weak chargeless and $SU(2)_{II}$ charged anti-quarks and anti-leptons. The anti-fermion states are reachable from the fermion states by the application of the discrete symmetry operator $\mathcal{C_N}\,\mathcal{P_N}$, presented in Ref. [38,39].

5 There are before the electroweak break two decoupled groups of four massless families of quarks and leptons, in the fundamental representations of $\widetilde{SU}(2)_{R,\widetilde{SO}(3,1)}\times\widetilde{SU}(2)_{II,\widetilde{SO}(4)}$ and $\widetilde{SU}(2)_{L,\widetilde{SO}(3,1)}\times\widetilde{SU}(2)_{I,\widetilde{SO}(4)}$ groups, respectively - the subgroups of $\widetilde{SO}(3,1)$ and $\widetilde{SO}(4)$ (Table 6.4, page 89). These eight families remain massless up to the electroweak break due to the "mass protection mechanism", that is due to the fact that the right handed members have no left handed partners with the same charges.

6 Before the electroweak break are all observable gauge fields massless: the gravity, the colour octet vector gauge fields (of the group $SU(3)$ from $SO(6)$), the weak triplet vector gauge field (of the group $SU(2)_I$ from $SO(4)$), and the hyper singlet vector gauge field (a superposition of $U(1)$ from $SO(6)$ and the third component of $SU(2)_{II}$ triplet). All are the superposition of the $f^{\alpha}{}_{c}\,\omega_{ab\alpha}$ spinor gauge fields.

7 There are scalar fields with the space index $(7,8)$ and with respect to the space index with the weak and the hyper charge of the Higgs's scalar. They belong with respect to

e. The starting action contains also the additional $SU(2)_{II}$ (from $SO(4)$) vector gauge triplet (one of the components contributes to the hyper charge gauge fiels as explained above in the footnote of **d**. of **C.i.**), as well as the scalar fields with the space index $s \in (5,6)$ and $t \in (9, 10, \ldots, 14)$. All these fields gain masses of the scale of the condensate (Table 7.1), which they interact with. They all are expressible with the superposition of $f^{\mu}{}_{m}\, \omega_{ab\mu}$ or of $f^{\mu}{}_{m}\, \tilde{\omega}_{ab\mu}$ [8].

C ii., C iii. There are many ways of breaking symmetries from $d = (13 + 1)$ to $d = (3 + 1)$. The assumed breaks explain the connection between the weak and the hyper charge and the handedness of spinors, manifesting correspondingly the observed properties of the family members - the quarks and the leptons, left and right handed (Table 6.2, page 86) - and of the observed vector gauge fields. After the break from $SO(13, 1)$ to $SO(3, 1) \times SU(2) \times U(1) \times SU(3)$ the anti-particles are accessible from particles by the application of the operator $\mathbb{C}_{\mathcal{N}} \cdot \mathcal{P}_{\mathcal{N}}$, as explained [9] in Refs. [38].

C iv. It is the condensate (Table 7.1) of two right handed neutrinos with the quantum numbers of one group of four families, which makes massive all the scalar gauge fields (those with the space index s equal to $(5, 6, 7, 8)$, as well as those with the space indexs equal to $(9, \ldots, 14)$) and those vector gauge fields, manifesting nonzero τ^4, τ^{23}, $\tilde{\tau}^4$, $\tilde{\tau}^{23}$, $\tilde{N}^3_R$ [2,1]. Only the vector gauge fields of Y, $SU(3)$ and $SU(2)$ remain massless, since they do not interact with the condensate.

C v. At the electroweak break the scalar fields with the space index $s = (7, 8)$ - originating in $\tilde{\omega}_{abs}$, as well as some superposition of $\omega_{s's"s}$ with the quantum numbers (Q, Q', Y'), conserving the colour and the electromagnetic charge - change their mutual interaction, and gaining nonzero vacuum expectation values change correspondingly also their masses. They contribute to mass matrices of twice the four families, as well as to the masses of the heavy vector bosons.

All the rest scalar fields keep masses of the scale of the condensate and are correspondingly unobservable in the low energy regime.

The fourth family to the observed three ones is *predicted to be observed* at the LHC. Its properties are under consideration [14,15], the *baryons of the stable family of the upper four families is offering the explanation for the dark matter* [13]. The

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additional quantum numbers either to one of the two groups of two triplets, (either to one of the two triplets of the groups $\widetilde{SU}(2)_{R\,\widetilde{SO}(3,1)}$ and $\widetilde{SU}(2)_{II\,\widetilde{SO}(4)}$, or to one of the two triplets of the groups $\widetilde{SU}(2)_{L\,\widetilde{SO}(3,1)}$ and $\widetilde{SU}(2)_{I\,\widetilde{SO}(4)}$, respectively), which couple through the family quantum numbers to one (the first two triplets) or to another (the second two triplets) group of four families - all are the superposition of $f^{\sigma}{}_{s}\ \tilde{\omega}_{ab\sigma}$, or they belong to three singlets, the scalar gauge fields of (Q, Q', Y'), which couple to the family members of both groups of families - they are the superposition of $f^{\sigma}{}_{s}\ \omega_{ab\sigma}$. Both kinds of scalar fields determine the fermion masses (Eq. (7.7)), offering the explanation for the higgs, the Yukawa couplings and the heavy bosons masses.

[8] In the case of free fields (if no spinor source, carrying their quantum numbers, is present) both $f^{\mu}{}_{m}\, \omega_{ab\mu}$ and $f^{\mu}{}_{m}\, \tilde{\omega}_{ab\mu}$ are expressible with vielbeins, correspondingly only one kind of the three gauge fields are the propagating fields.

[9] The discrete symmetry operator $\mathbb{C}_{\mathcal{N}} \cdot \mathcal{P}_{\mathcal{N}}$, Refs. [38,39], does not contain $\tilde{\gamma}^a$'s degrees of freedom. To each family member there corresponds the anti-member, with the same family quantum number.

triplet and anti-triplet scalar fields contribute together with the condensate to the *matter/antimatter assymetry*.

Let us (formally) rewrite that part of the action of Eq.(7.8), which determines the spinor degrees of freedom, in the way that we can clearly see that the action does in the low energy regime manifest by the *standard model* required degrees of freedom of the fermions, vector and scalar gauge fields [4,5,3,1,9,6–8,11–14].

$$\begin{aligned}\mathcal{L}_f = &\bar{\psi}\gamma^m (p_m - \sum_{A,i} g^A \tau^{Ai} A^{Ai}_m)\psi + \\ &\{\sum_{s=7,8} \bar{\psi}\gamma^s p_{0s}\, \psi\} + \\ &\{\sum_{t=5,6,9,\ldots,14} \bar{\psi}\gamma^t p_{0t}\, \psi\},\end{aligned} \tag{7.9}$$

where

$$\begin{aligned} p_{0s} &= p_s - \frac{1}{2} S^{s's"}\omega_{s's"s} - \frac{1}{2}\tilde{S}^{ab}\tilde{\omega}_{abs}, \\ p_{0t} &= p_t - \frac{1}{2} S^{t't"}\omega_{t't"t} - \frac{1}{2}\tilde{S}^{ab}\tilde{\omega}_{abt}, \end{aligned}$$

with $m \in (0,1,2,3)$, $s \in (7,8)$, $(s', s") \in (5,6,7,8)$, (a,b) (appearing in $\tilde{S}^{ab}$) run within either $(0,1,2,3)$ or $(5,6,7,8)$, t runs $\in (5,\ldots,14)$, $(t', t")$ run either $\in (5,6,7,8)$ or $\in (9,10,\ldots,14)$. The spinor function ψ represents all family members of all the $2^{\frac{7+1}{2}-1} = 8$ families.

The first line of Eq. (7.9) determines (in $d = (3+1)$) the kinematics and dynamics of spinor (fermion) fields, coupled to the vector gauge fields. The generators τ^{Ai} of the charge groups are expressible in terms of S^{ab} through the complex coefficients ${c^{Ai}}_{ab}$ [10],

$$\tau^{Ai} = \sum_{a,b} {c^{Ai}}_{ab}\, S^{ab}\,, \tag{7.10}$$

fulfilling the commutation relations

$$\{\tau^{Ai}, \tau^{Bj}\}_- = i\delta^{AB} f^{Aijk}\tau^{Ak}\,. \tag{7.11}$$

They represent the colour, the weak and the hyper charges (as well as the $SU(2)_{II}$ and τ^4 charges, the gauge fields of which gain masses interacting with the condensate, Table 7.1, leaving massless only the hyper charge vector gauge field). The corresponding vector gauge fields A^{Ai}_m are expressible with the spin connection fields ω_{stm}, with (s,t) either $\in (5,6,7,8)$ or $\in (9,\ldots,14)$, in agreement with the

[10] $\vec{\tau}^1 := \frac{1}{2}(S^{58} - S^{67}, S^{57} + S^{68}, S^{56} - S^{78})$, $\vec{\tau}^2 := \frac{1}{2}(S^{58} + S^{67}, S^{57} - S^{68}, S^{56} + S^{78})$, $\vec{\tau}^3 := \frac{1}{2}\{S^{9\,12} - S^{10\,11}, S^{9\,11} + S^{10\,12}, S^{9\,10} - S^{11\,12}, S^{9\,14} - S^{10\,13}, S^{9\,13} + S^{10\,14}, S^{11\,14} - S^{12\,13}, S^{11\,13} + S^{12\,14}, \frac{1}{\sqrt{3}}(S^{9\,10} + S^{11\,12} - 2S^{13\,14})\}$, $\tau^4 := -\frac{1}{3}(S^{9\,10} + S^{11\,12} + S^{13\,14})$. After the electroweak break the charges $Y := \tau^4 + \tau^{23}$, $Y' := -\tau^4 \tan^2\vartheta_2 + \tau^{23}$, $Q := \tau^{13} + Y$, $Q' := -Y\tan^2\vartheta_1 + \tau^{13}$ manifest. θ_1 is the electroweak angle, breaking $SU(2)_I$, θ_2 is the angle of the break of the $SU(2)_{II}$ from $SU(2)_I \times SU(2)_{II}$.

assumptions **A ii.** and **A iii.**. I demonstrate in Ref. [1] the equivalence between the usual Kaluza-Klein procedure leading to the vector gauge fields through the vielbeins and the procedure with the spin connections proposed by the *spin-charge-family* theory.

All vector gauge fields, appearing in the first line of Eq. (7.9), except $A^{2\pm}_m$ and $A^{Y'}_m$ ($= \cos\vartheta_2\, A^{23}_m - \sin\vartheta_2\, A^4_m$, Y' and τ^4 are defined in [11]), are massless before the electroweak break. $\vec{A}^3_m$ carries the colour charge $SU(3)$ (originating in $SO(6)$), $\vec{A}^1_m$ carries the weak charge $SU(2)_I$ ($SU(2)_I$ and $SU(2)_{II}$ are the subgroups of $SO(4)$) and A^Y_m ($= \sin\vartheta_2\, A^{23}_m + \cos\vartheta_2\, A^4_m$) carries the corresponding $U(1)$ charge ($Y = \tau^{23} + \tau^4$, τ^4 originates in $SO(6)$ and τ^{23} is the third component of the second $SU(2)_{II}$ group, A^4_m and $\vec{A}^2_m$ are the corresponding vector gauge fields). The fields $A^{2\pm}_m$ and $A^{Y'}_m$ get masses of the order of the condensate scale through the interaction with the condensate of the two right handed neutrinos with the quantum numbers of one - the upper one - of the two groups of four families (the assumption **A iv.**, Table 7.1). (See Ref. [1].)

Since spinors (fermions) carry besides the family members quantum numbers also the family quantum numbers, determined by $\tilde{S}^{ab} = \frac{i}{4}(\tilde{\gamma}^a\tilde{\gamma}^b - \tilde{\gamma}^b\tilde{\gamma}^a)$, there are correspondingly $2^{(7+1)/2-1} = 8$ families [1], which split into two groups of families, each manifesting the $(\widetilde{SU}(2)_{\widetilde{SO}(3,1)} \times \widetilde{SU}(2)_{\widetilde{SO}(4)} \times U(1))$ symmetry.

If there are no fermions present then the vector gauge fields of the family members and of the family charges - ω_{abm} and $\tilde{\omega}_{abm}$, respectively - are uniquely expressible with the vielbeins [2,1].

The scalar fields, the gauge fields with the space index $s = (7, 8)$, which are either superposition of $\tilde{\omega}_{abs}$ or of $\omega_{s'ts}$, determine - after gaining nonzero vacuum expectation values (the assumption **A v.** and comments **C v.**) - masses of fermions (belonging to two groups of four families of family members of spinors) and weak bosons.

The condensate (the assumption **A iv.**), Table 7.1, gives masses of the order of the scale of its appearance to all the scalar gauge fields, presented in the second and the third line of Eq. (7.9).

The vector gauge fields of the (before the electroweak break) conserved the colour, the weak and the hyper charges ($\vec{\tau}^3$, $\vec{\tau}^1$, Y) do not interact with the condensate and stay correspondingly massless. After the electroweak break - when the scalar fields (those with the family quantum numbers and those with the family members quantum numbers (Q, Q', Y')) with the space index $s = (7, 8)$ start to strongly self interact, gaining nonzero vacuum expectation values - only the charges $\vec{\tau}^3$ and $Q = Y + \tau^{13}$ are the conserved charges. No family quantum numbers are conserved, since all the scalar fields with the family quantum numbers and the space index $s = (7, 8)$ gain nonzero vacuum expectation values.

Quarks and leptons have the "spinor" quantum number (τ^4, originating in $SO(6)$), presented in Table 6.2, page 86) equal to $\frac{1}{6}$ and $-\frac{1}{2}$, respectively.

[11] $Y' := -\tau^4 \tan^2\vartheta_2 + \tau^{23}, \tau^4 = -\frac{1}{3}(S^{9\,10} + S^{11\,12} + S^{13\,14})$.

state	S^{03}	S^{12}	τ^{13}	τ^{23}	τ^{4}	Y	Q	$\tilde{\tau}^{13}$	$\tilde{\tau}^{23}$	$\tilde{\tau}^{4}$	$\tilde{Y}$	$\tilde{Q}$	$\tilde{N}^3_L$	$\tilde{N}^3_R$
$(\|\nu^{\mathbf{VIII}}_{\mathbf{1R}} >_{\mathbf{1}} \|\nu^{\mathbf{VIII}}_{\mathbf{2R}} >_{\mathbf{2}})$	0	0	0	1	−1	0	0	0	1	−1	0	0	0	1
$(\|\nu^{VIII}_{1R} >_1 \|e^{VIII}_{2R} >_2)$	0	0	0	0	−1	−1	−1	0	1	−1	0	0	0	1
$(\|e^{VIII}_{1R} >_1 \|e^{VIII}_{2R} >_2)$	0	0	0	−1	−1	−2	−2	0	1	−1	0	0	0	1

Table 7.1. This table is taken from [2]. The condensate of the two right handed neutrinos ν_R, with the VIII^{th} family quantum numbers, coupled to spin zero and belonging to a triplet with respect to the generators τ^{2i}, is presented, together with its two partners. The right handed neutrino has $Q = 0 = Y$. The triplet carries $\tau^4 = -1$, $\tilde{\tau}^{23} = 1$, $\tilde{\tau}^4 = -1$, $\tilde{N}^3_R = 1$, $\tilde{N}^3_L = 0$, $\tilde{Y} = 0$, $\tilde{Q} = 0$. The family quantum numbers are presented in Table 6.4, page 89.

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