Can the matter-antimatter asymmetry be easier to understand within the "spin-charge-family-theory", predicting twice four families and two times $SU(2)$ vector gauge and scalar fields?

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Abstract

This contribution is an attempt to try to understand the matter-antimatter asymmetry in the universe within the *spin-charge-family-theory* [1, 2] if assuming that transitions in non equilibrium processes among instanton vacua and complex phases in mixing matrices are the sources of the matter-antimatter asymmetry, as studied in the literature [3–6] for several proposed theories. The *spin-charge-family-theory* is, namely, very promising in showing the right way beyond the *standard model*. It predicts families and their mass matrices, explaining the origin of the charges and of the gauge fields. It predicts that there are, after the universe passes through two $SU(2)\times U(1)$ phase transitions, in which the symmetry breaks from $SO(1,3)\times SU(2)\times SU(2)\times U(1)\times SU(3)$ first to $SO(1,3)\times SU(2)\times U(1)\times SU(3)$ and then to $SO(1,3)\times U(1)\times SU(3)$, twice decoupled four families. The upper four families gain masses in the first phase transition, while the second four families gain masses at the electroweak break. To these two breaks of symmetries the scalar non Abelian fields, the (superposition of the) gauge fields of the operators generating families, contribute. The lightest of the upper four families is stable (in comparison with the life of the universe) and is therefore a candidate for constituting the dark matter. The heaviest of the lower four families should be seen at the LHC or at somewhat higher energies.

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I. INTRODUCTION

The *theory unifying spin and charges and predicting families* (*spin-charge-family-theory*) assumes that spinors carry in $d \geq 4$ ($d = 1 + 13$ is studied) only two kinds of the spin. The Dirac kind γ^a manifests after several appropriate breaks of the starting symmetry as the spin and all the charges. The second kind called $\{\gamma^a$ ($\{\gamma^a, \tilde{\gamma}^b\}_+ = 0$) generates families. Accordingly there are in $d \geq 4$, besides the vielbeins, also the two kinds of the spin connection fields, which are the gauge fields of the corresponding operators S^{ab} and $\tilde{S}^{ab}$. Those connected with S^{ab} manifest in $d = (1 + 3)$ as the vector gauge fields, while those connected with $\tilde{S}^{ab}$ manifest as the scalar fields and determine on the tree level the mass matrices.

Let me make a short review of the so far made predictions of the *spin-charge-family-theory*:

- The *spin-charge-family-theory* has the explanation for the appearance of the internal degrees of freedom – the spin and the charges while unifying them under the assumption that the universe went through several phase transitions which cause the appropriate breaks of the starting symmetry. Then the fact that the right handed (with respect to SO(1,3)) fermions are weak chargeless, while the left handed ones carry the weak charge emerges, as well as that there exist leptons (singlets with respect to the colour charge) and quarks (triplets with respect to the colour charge) [1, 2].

- The theory explains the appearance of massless families at the low energy regime under the assumption that there are breaks which leave the massless fermions of only one handedness [7]. Assuming that breaks of symmetries affect the whole internal space — the space defined by both kinds of the Clifford algebra objects — it predicts in the energy regime close below 10^{16} GeV eight massless families. The manifested symmetry is (assumed to be) at this stage $SO(1,3) \times SO(4) \times U(1) \times SU(3)$. The next break of the symmetry of the universe to $SO(1,3) \times SU(2) \times U(1) \times SU(3)$ leaves four families massless [2], while the vacuum expectation values of superposition of the starting fields which manifest in $(1+3)$ as scalar fields, make the upper four families and the corresponding gauge fields massive. After the electroweak break also the lower four families become massive due to the vacuum expectation values of superposition of the starting fields, together with the weak bosons.

- The theory predicts the fourth family, which will be observed at the LHC or at somewhat higher energies [8], and the fifth stable family (with no mixing matrix elements couplings to the lower four families in comparison with the age of the universe), the baryons and neutrinos of which are the candidates to form the dark matter.

- The masses of this fifth family members are according to the so far made rough estimations [2, 8] larger than a few TeV and smaller than 10^{10} TeV. The members of the family have approximately the same mass, at least on the tree level [9].

- The studies [10] of the history of the stable fifth family members in the evolution of the universe and of their interactions with the ordinary matter in the DAMAs and the CDMSs experiments done so far lead to the prediction that the masses of the fifth family members, if they constitute the dark matter, are a few hundred TeV, independent of the fifth family fermion-antifermion asymmetry. The Xe experiment looks like to be in disagreement, but careful analyses show that one should wait for further data [11] to make the final conclusion.

  The lightest fifth family baryon is, in the case that all the quarks have approximately (within a hundred GeV) the same mass [10], the fifth family neutron, due to the attractive electromagnetic interaction. The difference in the weak interaction can be for large enough masses neglected.

- The fermion asymmetry in the *approach* has not yet really been studied.

- The studies [10] of the evolution of this stable fifth family members rely on my rough estimations [10] of the behaviour of the coloured fifth family objects (single quarks and antiquarks or coloured pairs of quarks or of antiquarks) during the colour phase transition. These estimations namely suggest that the coloured objects either annihilate with the anti-objects or they form colourless neutrons and antineutrons and correspondingly decouple from the plasma soon after the colour phase transition starts, due to the very strong binding energy of the fifth family baryons (with respect to the first family baryons) long enough before the first family quarks start to form the baryons. These estimations should be followed by more accurate studies.

- The so far done studiessuggest strongly that the number density of the fifth family neutrinos (of approximately the same mass as the fifth family quarks and leptons), which also contribute to the dark matter, is pretty much reduced due to the neutrino-antineutrino annihilation closed below the electroweak break. The weak annihilation cross section is expected to play much stronger role for neutrinos than for strongly bound fifth family quarks in the fifth family neutron (due to the huge binding energy of the fifth family quarks), what also remains to be proved.

- The estimations [8] of the properties of the lower four families on the tree level call for the calculations beyond the tree level, which should hopefully demonstrate, that the loop corrections (in all orders) bring the main differences in the properties of the family members. These calculations are in progress [12].

Although we can say that the *spin-charge-family-theory* looks very promising as the right way to explain where do the assumptions of the *standard model* originate, there are obviously many not yet studied, or at least far from being carefully enough studied open problems. Many a problem is common to all the theories, like the first family baryon asymmetry, which I am going to discuss within the *spin-charge-family-theory* in this contribution. Some of the problems are common to all the theories assuming more than so far observed $(1+3)$ dimensions, in particular the *spin-charge-family-theory* shares some problems with all the Kaluza-Klein-like theories. We are trying to solve them first on toy models [7].

The main new step in the *spin-charge-family-theory* — the explanation of the appearance of families by assuming that both existing kinds of the Clifford algebra objects should be used to treat correctly the fermion degrees of freedom — limits very much the properties of families and their members. The simple starting action in $d = (1+13)$, which in $d = (1+3)$ demonstrates the mass matrices, namely fixes to high extent the fermion properties after the breaks of symmetries. Therefore this proposal might soon be studied accurately enough to show whether it is the right theory or not.

This contribution is an attempt to try to understand what can the *spin-charge-family-theory* say about the fermion-antifermion asymmetry when taking into account the proposals of the references [4–6] (and of the works cited therein). These works study the soliton solutions of non Abelian gauge fields with many different vacua and evaluate fermion number nonconservation due to possible transitions among different vacua in non equilibrium processes during the phase transitions through which the universe passed. In such processes fermion (and also antifermion) currents are not conserved since CP is not nonconserved. To the CP nonconservation also the complex matrix elements determining the transitions among families contribute and consequently influence the first family fermion-antifermion asymmetry.

Since the *spin-charge-family-theory* predicts below the unification scale of all the charges two kinds of phase transitions (first from $SO(1,3) \times SU(2) \times SU(2) \times U(1) \times SU(3)$ to $SO(1,3) \times SU(2) \times U(1) \times SU(3)$, in which the upper four families gain masses and so do the corresponding vector gauge fields, and then from $SO(1,3) \times SU(2) \times U(1) \times SU(3)$ to $SO(1,3) \times U(1) \times SU(3)$, in which the lower four families and the corresponding gauge fields gain masses), in which besides the vector gauge fields also the scalar gauge fields (the gauge fields of $\tilde{S}^{ab}$ and also of S^{ab} with the scalar index with respect to $(1+3)$) contribute, the fermion-antifermion asymmetry might very probably have for the stable fifth family an opposite sign than for the first family.

It might therefore be that the existence of two kinds of four families, together with two kinds of the vector gauge fields and two kinds of the scalar fields help to easier understanding the first family fermion-antifermion asymmetry.

Although I am studying the fermion asymmetry, together with the discrete symmetries, in the *spin-charge-family-theory* for quite some time (not really intensively), this contribution is stimulated by the question of M.Y. Khlopov [13], since he is proposing the scenario, in which my stable fifth family members should manifest an opposite fermion asymmetry than the first family members, that is antifermion-fermion asymmetry. While in the case that the fifth family members have masses around 100 TeV or higher and the neutron is the lightest baryon and neutrino the lightest lepton [10] the fifth family baryon asymmetry plays no role (since in this case the fifth family neutrons and neutrinos as well as their antiparticles interact weakly enough among themselves and with the ordinary matter that the assumption that they constitute the dark matter is in agreement with the observations). Maxim [14] claims that the fifth family members with the quark masses not higher than 10 TeV are also the candidates for the dark matter, provided that $\bar{u}_5\bar{u}_5\bar{u}_5$ is the lightest antibaryon and that there is an excess of antibaryons over the baryons in the fifth family case.

II. A SHORT OVERVIEW OF THE THEORY UNIFYING SPIN AND CHARGES AND EXPLAINING FAMILIES

In this section I briefly repeat the main ideas of the *spin-charge-family-theory*. I kindly ask the reader to learn more about this theory in the references [1, 2] as well as in my talk presented in this proceedings and in the references therein.

I am proposing a simple action in $d = (1 + 13)$-dimensional space. Spinors carry two kinds of the spin (no charges).
i. The Dirac spin, described by γ^a's, defines the spinor representation in $d = (1+13)$. After the break of the starting symmetry $SO(1,13)$ (through $SO(1,7) \times SO(6)$) to the symmetry of the *standard model* in $d = (1+3)$ ($SO(1,3) \times U(1) \times SU(2) \times SU(3)$) it defines the hyper charge ($U(1)$), the weak charge ($SU(2)$, with the left handed representation of $SO(1,3)$ manifesting naturally the weak charge and the right handed ones appearing as the weak singlets) and the colour charge ($SU(3)$).
ii. The second kind of the spin [1], described by $\tilde{\gamma}^a$'s ($\{\tilde{\gamma}^a, \tilde{\gamma}^b\}_+ = 2\,\eta^{ab}$) and anticommuting with the Dirac γ^a ($\{\gamma^a, \tilde{\gamma}^b\}_+ = 0$), defines the families of spinors.
Accordingly spinors interact with the two kinds of the spin connection fields and the vielbeins.

We have

$$\begin{aligned}
&\{\gamma^a, \gamma^b\}_+ = 2\eta^{ab} = \{\tilde{\gamma}^a, \tilde{\gamma}^b\}_+, \quad \{\gamma^a, \tilde{\gamma}^b\}_+ = 0,\\
&S^{ab} := (i/4)(\gamma^a\gamma^b - \gamma^b\gamma^a), \quad \tilde{S}^{ab} := (i/4)(\tilde{\gamma}^a\tilde{\gamma}^b - \tilde{\gamma}^b\tilde{\gamma}^a), \quad \{S^{ab}, \tilde{S}^{cd}\}_- = 0.
\end{aligned} \tag{1}$$

The action

$$\begin{aligned}
S \;&=\; \int \; d^dx \; E \; \mathcal{L}_f + \\
&\quad\;\; \int \; d^dx \; E \; (\alpha \, R + \tilde{\alpha} \, \tilde{R}) \, ,
\end{aligned} \tag{2}$$

$$\begin{aligned}
\mathcal{L}_f \;&=\; \frac{1}{2} \, (E\bar{\psi} \, \gamma^a p_{0a} \psi) + h.c. \, ,\\
p_{0a} \;&=\; f^\alpha{}_a p_{0\alpha} + \frac{1}{2E} \, \{p_\alpha, E f^\alpha{}_a\}_- ,
\end{aligned}$$

$$
\begin{aligned}
p_{0\alpha} &= p_\alpha - \frac{1}{2} S^{ab}\omega_{ab\alpha} - \frac{1}{2}\tilde{S}^{ab}\tilde{\omega}_{ab\alpha}\,,\\
R &= \frac{1}{2}\,\{f^{\alpha[a}f^{\beta b]}\;(\omega_{ab\alpha,\beta} - \omega_{ca\alpha}\,\omega^{c}{}_{b\beta})\} + h.c.\,,\\
\tilde{R} &= \frac{1}{2}\,f^{\alpha[a}f^{\beta b]}\;(\tilde{\omega}_{ab\alpha,\beta} - \tilde{\omega}_{ca\alpha}\tilde{\omega}^{c}{}_{b\beta}) + h.c.\,,
\end{aligned} \tag{3}
$$

manifests ($f^{\alpha[a}f^{\beta b]} = f^{\alpha a}f^{\beta b} - f^{\alpha b}f^{\beta a}$) after the break of symmetries all the known gauge fields and the scalar fields, and the mass matrices. To see the manifestation of the covariant momentum and the mass matrices we rewrite formally the action for a Weyl spinor in $d = (1+13)$ as follows

$$
\begin{aligned}
\mathcal{L}_f &= \bar{\psi}\gamma^m(p_m - \sum_{A,i} g^A\tau^{Ai}A^{Ai}_m)\psi + \\
&\quad \{\sum_{s=7,8} \bar{\psi}\gamma^s p_{0s}\,\psi\} + \\
&\quad \text{the rest},
\end{aligned} \tag{4}
$$

where $m = 0,1,2,3$ with

$$
\begin{aligned}
\tau^{Ai} &= \sum_{a,b} c^{Ai}{}_{ab}\,S^{ab},\\
\{\tau^{Ai},\tau^{Bj}\}_- &= i\delta^{AB}f^{Aijk}\tau^{Ak}.
\end{aligned} \tag{5}
$$

All the charges and the spin of one family are determined by S^{ab}, with S^{ab} as the only internal degree of freedom of one family (besides the family quantum number, determined by $\tilde{S}^{ab}$), manifesting after the breaks at the low energy regime as the generators of the observed groups (Eq. (5)) $U(1), SU(2)$ and $SU(3)$, for $A = 1,2,3$, respectively.

The breaks of the starting symmetry from $SO(1,13)$ to the symmetry $SO(1,7)\times SU(3)\times U(1)$ and further to $SO(1,3)\times SU(2)\times SU(2)\times U(1)\times SU(3)$ are assumed to leave all the low lying families of spinors massless. There are eight such massless families ($2^{8/2-1}$) before further breaks.

Accordingly the first row of the action in Eq. (4) manifests the effective *standard model* fermions part of the action before the weak break, while the second part manifests, after the appropriate breaks of symmetries (when $\omega_{ab\sigma}$ and $\tilde{\omega}_{ab\sigma}$, $\sigma \in (5,6,7,8)$, fields gain the nonzero vacuum expectation values on the tree level) the mass matrices.

The generators $\tilde{S}^{ab}$ take care of the families, transforming each member of one family into the corresponding member of another family, due to the fact that $\{S^{ab}, \tilde{S}^{cd}\}_- = 0$ (Eq.(1)).

i		$\vert^a\psi_i>$	$\Gamma^{(1,3)}$	S^{12}	$\Gamma^{(4)}$	τ^{13}	τ^{23}	Y	Q
		Octet, $\Gamma^{(1,7)} = 1,\ \Gamma^{(6)} = -1,$ of quarks							
1	u_R^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{(+)}\vert\stackrel{56}{(+)}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	1	$\frac{1}{2}$	1	0	$\frac{1}{2}$	$\frac{2}{3}$	$\frac{2}{3}$
2	u_R^{c1}	$\stackrel{03}{[-i]}\stackrel{12}{[-]}\vert\stackrel{56}{(+)}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	1	$-\frac{1}{2}$	1	0	$\frac{1}{2}$	$\frac{2}{3}$	$\frac{2}{3}$
3	d_R^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{(+)}\vert\stackrel{56}{[-]}\stackrel{78}{[-]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	1	$\frac{1}{2}$	1	0	$-\frac{1}{2}$	$-\frac{1}{3}$	$-\frac{1}{3}$
4	d_R^{c1}	$\stackrel{03}{[-i]}\stackrel{12}{[-]}\vert\stackrel{56}{[-]}\stackrel{78}{[-]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	1	$-\frac{1}{2}$	1	0	$-\frac{1}{2}$	$-\frac{1}{3}$	$-\frac{1}{3}$
5	d_L^{c1}	$\stackrel{03}{[-i]}\stackrel{12}{(+)}\vert\stackrel{56}{[-]}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	-1	$\frac{1}{2}$	-1	$-\frac{1}{2}$	0	$\frac{1}{6}$	$-\frac{1}{3}$
6	d_L^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{[-]}\vert\stackrel{56}{[-]}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	-1	$-\frac{1}{2}$	-1	$-\frac{1}{2}$	0	$\frac{1}{6}$	$-\frac{1}{3}$
7	u_L^{c1}	$\stackrel{03}{[-i]}\stackrel{12}{(+)}\vert\stackrel{56}{(+)}\stackrel{78}{[-]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	-1	$\frac{1}{2}$	-1	$\frac{1}{2}$	0	$\frac{1}{6}$	$\frac{2}{3}$
8	u_L^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{[-]}\vert\stackrel{56}{(+)}\stackrel{78}{[-]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	-1	$-\frac{1}{2}$	-1	$\frac{1}{2}$	0	$\frac{1}{6}$	$\frac{2}{3}$

TABLE I: The 8-plet of quarks - the members of $SO(1,7)$ subgroup of the group $SO(1,13)$, belonging to one Weyl left handed ($\Gamma^{(1,13)} = -1 = \Gamma^{(1,7)} \times \Gamma^{(6)}$) spinor representation of $SO(1,13)$. It contains the left handed weak charged quarks and the right handed weak chargeless quarks of a particular colour $(1/2, 1/(2\sqrt{3}))$. Here $\Gamma^{(1,3)}$ defines the handedness in $(1+3)$ space, S^{12} defines the ordinary spin (which can also be read directly from the basic vector, both vectors with both spins, $\pm\frac{1}{2}$, are presented), τ^{13} defines the third component of the weak charge, τ^{23} the third component of the $SU(2)_{II}$ charge, τ^4 (the $U(1)$ charge) defines together with the τ^{23} the hyper charge ($Y = \tau^4 + \tau^{23}$), $Q = Y + \tau^{13}$ is the electromagnetic charge. The reader can find the whole Weyl representation in the ref. [16].

Using the technique [15] and analysing the vectors as the eigenvectors of the *standard model* groups we present vectors in the space of charges and spins in terms of projectors and nilpotents as can be learned in Appendix, in the references [1, 2] and also in my talk in the Proceedings of Bled workshop 2010.

I present in Table I the eightplet (the representation of $SO(1,7)$ of quarks of a particular colour charge ($\tau^{33} = 1/2$, $\tau^{38} = 1/(2\sqrt{3})$), and $U(1)$ charge ($\tau^4 = 1/6$) and on Table II the eightplet of the corresponding (colour chargeless) leptons.

In both tables the vectors are chosen to be the eigenvectors of the operators of handedness

i		$\vert^{a}\psi_i>$	$\Gamma^{(1,3)}$	S^{12}	$\Gamma^{(4)}$	τ^{13}	τ^{23}	Y	Q
		Octet, $\Gamma^{(1,7)}=1,\ \Gamma^{(6)}=-1$, of quarks							
1	ν_R	$\stackrel{03}{(+i)}\stackrel{12}{(+)} \vert \stackrel{56}{(+)}\stackrel{78}{(+)} \vert\vert \stackrel{9\,10}{(+)}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	1	$\frac{1}{2}$	1	0	$\frac{1}{2}$	0	0
2	ν_R	$\stackrel{03}{[-i]}\stackrel{12}{[-]} \vert \stackrel{56}{(+)}\stackrel{78}{(+)} \vert\vert \stackrel{9\,10}{(+)}\ \stackrel{11\,12}{[-]}\ \stackrel{13\,14}{[-]}$	1	$-\frac{1}{2}$	1	0	$\frac{1}{2}$	0	0
3	e_R	$\stackrel{03}{(+i)}\stackrel{12}{(+)} \vert \stackrel{56}{[-]}\stackrel{78}{[-]} \vert\vert \stackrel{9\,10}{(+)}\ \stackrel{11\,12}{[-]}\ \stackrel{13\,14}{[-]}$	1	$\frac{1}{2}$	1	0	$-\frac{1}{2}$	-1	-1
4	e_R	$\stackrel{03}{[-i]}\stackrel{12}{[-]} \vert \stackrel{56}{[-]}\stackrel{78}{[-]} \vert\vert \stackrel{9\,10}{(+)}\ \stackrel{11\,12}{[-]}\ \stackrel{13\,14}{[-]}$	1	$-\frac{1}{2}$	1	0	$-\frac{1}{2}$	-1	-1
5	e_L	$\stackrel{03}{[-i]}\stackrel{12}{(+)} \vert \stackrel{56}{[-]}\stackrel{78}{(+)} \vert\vert \stackrel{9\,10}{(+)}\ \stackrel{11\,12}{[-]}\ \stackrel{13\,14}{[-]}$	-1	$\frac{1}{2}$	-1	$-\frac{1}{2}$	0	$-\frac{1}{2}$	-1
6	e_L	$\stackrel{03}{(+i)}\stackrel{12}{[-]} \vert \stackrel{56}{[-]}\stackrel{78}{(+)} \vert\vert \stackrel{9\,10}{(+)}\ \stackrel{11\,12}{[-]}\ \stackrel{13\,14}{[-]}$	-1	$-\frac{1}{2}$	-1	$-\frac{1}{2}$	0	$-\frac{1}{2}$	-1
7	ν_L	$\stackrel{03}{[-i]}\stackrel{12}{(+)} \vert \stackrel{56}{(+)}\stackrel{78}{[-]} \vert\vert \stackrel{9\,10}{(+)}\ \stackrel{11\,12}{[-]}\ \stackrel{13\,14}{[-]}$	-1	$\frac{1}{2}$	-1	$\frac{1}{2}$	0	$-\frac{1}{2}$	0
8	ν_L	$\stackrel{03}{(+i)}\stackrel{12}{[-]} \vert \stackrel{56}{(+)}\stackrel{78}{[-]} \vert\vert \stackrel{9\,10}{(+)}\ \stackrel{11\,12}{[-]}\ \stackrel{13\,14}{[-]}$	-1	$-\frac{1}{2}$	-1	$\frac{1}{2}$	0	$-\frac{1}{2}$	0

TABLE II: The 8-plet of leptons - the members of $SO(1,7)$ subgroup of the group $SO(1,13)$, belonging to one Weyl left handed ($\Gamma^{(1,13)} = -1 = \Gamma^{(1,7)} \times \Gamma^{(6)}$) spinor representation of $SO(1,13)$. It contains the colour chargeless left handed weak charged leptons and the right handed weak chargeless leptons. The rest of notation is the same as in Table II.

$\Gamma^{(n)}$, the generators τ^{13}, τ^{23}, τ^{33} τ^{38}, $Y = \tau^4 + \tau^{23}$ and $Q = Y + \tau^{13}$. They are also eigenvectors of the corresponding $\tilde{S}^{ab}$, $\tilde{\tau}^{Ai}, A = 1,2,3$ and $\tilde{Y}, \tilde{Q}$. One easily sees that the right handed vectors (with respect to $SO(1,3)$) are weak ($SU(2)_I$) chargeless and are doublets with respect to the second $SU(2)_{II}$, while the left handed are weak charged and singlets with respect to $SU(2)_{II}$.

The generators $\tilde{S}^{ab}$ transform each member of a family into the same member of other $2^{\frac{8}{2}-1}$ families. The eight families of the first member of the eightplet of quarks from Table I, for example, that is of the right handed u-quark of the spin $\frac{1}{2}$, are presented in the left column of Table III. The corresponding right handed neutrinos, belonging to eight different families, are presented in the right column of the same table. The u-quark member of the eight families and the ν members of the same eight families are generated by $\tilde{S}^{cd}$, $c, d \in \{0,1,2,3,5,6,7,8\}$ from any starting family.

Let us present also the quantum numbers of the families from Table III. In Table IV the handedness of the families $\tilde{\Gamma}^{(1+3)}(= -4i\tilde{S}^{03}\tilde{S}^{12})$, $\tilde{S}^{03}_L, \tilde{S}^{12}_L$, $\tilde{S}^{03}_R, \tilde{S}^{12}_R$ (the diagonal matrices of

I_R	u_R^{c1}	$\stackrel{03}{[+i]}\stackrel{12}{(+)}\vert\stackrel{56}{(+)}\stackrel{78}{[+]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	ν_R	$\stackrel{03}{[+i]}\stackrel{12}{(+)}\vert\stackrel{56}{(+)}\stackrel{78}{[+]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$
II_R	u_R^{c1}	$\stackrel{03}{[+i]}\stackrel{12}{(+)}\vert\stackrel{56}{[+]}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	ν_R	$\stackrel{03}{(+i)}\stackrel{12}{[+]}\vert\stackrel{56}{(+)}\stackrel{78}{[+]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$
III_R	u_R^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{[+]}\vert\stackrel{56}{(+)}\stackrel{78}{[+]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	ν_R	$\stackrel{03}{(+i)}\stackrel{12}{[+]}\vert\stackrel{56}{[+]}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$
IV_R	u_R^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{[+]}\vert\stackrel{56}{[+]}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	ν_R	$\stackrel{03}{[+i]}\stackrel{12}{(+)}\vert\stackrel{56}{[+]}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$
V_R	u_R^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{(+)}\vert\stackrel{56}{(+)}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	ν_R	$\stackrel{03}{(+i)}\stackrel{12}{(+)}\vert\stackrel{56}{(+)}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$
VI_R	u_R^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{(+)}\vert\stackrel{56}{[+]}\stackrel{78}{[+]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	ν_R	$\stackrel{03}{(+i)}\stackrel{12}{(+)}\vert\stackrel{56}{[+]}\stackrel{78}{[+]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$
VII_R	u_R^{c1}	$\stackrel{03}{[+i]}\stackrel{12}{[+]}\vert\stackrel{56}{(+)}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	ν_R	$\stackrel{03}{[+i]}\stackrel{12}{[+]}\vert\stackrel{56}{(+)}\stackrel{78}{(+)}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$
$VIII_R$	u_R^{c1}	$\stackrel{03}{[+i]}\stackrel{12}{[+]}\vert\stackrel{56}{[+]}\stackrel{78}{[+]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{[-]}\;\stackrel{13\,14}{[-]}$	ν_R	$\stackrel{03}{[+i]}\stackrel{12}{[+]}\vert\stackrel{56}{[+]}\stackrel{78}{[+]}\Vert\stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$

TABLE III: Eight families of the right handed u_R quark with the spin $\frac{1}{2}$, the colour charge $\tau^{33} = 1/2$, $\tau^{38} = 1/(2\sqrt{3})$ and of the colourless right handed neutrino ν_R of the spin $\frac{1}{2}$ are presented in the left and in the right column, respectively. $S^{ab}, a, b \in \{0,1,2,3,5,6,7,8\}$ transform u_R^{c1} of the spin $\frac{1}{2}$ and the chosen colour $c1$ to all the members of the same colour: to the right handed u_R^{c1} of the spin $-\frac{1}{2}$, to the left u_L^{c1} of both spins $(\pm\frac{1}{2})$, to the right handed d_R^{c1} of both spins $(\pm\frac{1}{2})$ and to the left handed d_L^{c1} of both spins $(\pm\frac{1}{2})$. They transform equivalently the right handed neutrino ν_R of the spin $\frac{1}{2}$ to the right handed ν_R of the spin $(-\frac{1}{2})$, to ν_L of both spins, to e_R of both spins and to e_L of both spins. $\tilde{S}^{ab}, a, b \in \{0,1,2,3,5,6,7,8\}$ transform a chosen member of one family into the same member of all the eight families.

$SO(1,3)$), $\tilde{\tau}^{13}$ (of one of the two $SU(2)_I$), $\tilde{\tau}^{23}$ (of the second $SU(2)_{II}$) are presented.

We see in Table IV that four of the eight families are singlets with respect to one of the two $SU(2)$ ($SU(2)_I$) groups determined by $\tilde{S}^{ab}$ and doublets with respect to the second $SU(2)$ ($SU(2)_{II}$), while the remaining four families are doublets with respect to the first $SU(2)_I$ and singlets with respect to the second $SU(2)_{II}$. When the first break appears, to which besides the vielbeins also the spin connections contribute, we expect that if only one of the two $SU(2)$ subgroups of $SO(1,7) \times U(1)$ breaking into $SO(1,3) \times SU(2) \times U(1)$ contributes in the break [9], namely that of the charges $\tilde{\tau}^{2i}$, together with $\tilde{N}^i_-$, there will be four families massless and mass protected after this break, namely those, which are singlets with respect to $\vec{\tilde{\tau}}^2$ and with respect to $\tilde{N}^i_-$ (Table IV), while for the other four families the vacuum expectation values of the scalars (particular combinations of vielbeins $f^{\sigma}{}_s$, and spin connections $\tilde{\omega}_{abs}, s \in \{5,8\}$) will take care of the mass matrices on the tree level and beyond

i	$\tilde{\Gamma}^{(1+3)}$	$\tilde{S}^{03}_L$	$\tilde{S}^{12}_L$	$\tilde{S}^{03}_R$	$\tilde{S}^{12}_R$	$\tilde{\tau}^{13}$	$\tilde{\tau}^{23}$	$\tilde{\tau}^{4}$	$\tilde{Y}'$	$\tilde{Y}$	$\tilde{Q}$
1	-1	$-\frac{i}{2}$	$\frac{1}{2}$	0	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	0
2	-1	$-\frac{i}{2}$	$\frac{1}{2}$	0	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	-1
3	-1	$\frac{i}{2}$	$-\frac{1}{2}$	0	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	0
4	-1	$\frac{i}{2}$	$-\frac{1}{2}$	0	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	-1
5	1	0	0	$\frac{i}{2}$	$\frac{1}{2}$	0	$\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	0	0
6	1	0	0	$\frac{i}{2}$	$\frac{1}{2}$	0	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	-1	-1
7	1	0	0	$-\frac{i}{2}$	$-\frac{1}{2}$	0	$\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	0	0
8	1	0	0	$-\frac{i}{2}$	$-\frac{1}{2}$	0	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	-1	-1

TABLE IV: The quantum numbers of each member of the eight families presented in Table III are presented: The handedness of the families $\tilde{\Gamma}^{(1+3)} = -4i\tilde{S}^{03}\tilde{S}^{12}$, the left and right handed $SO(1,3)$ quantum numbers $\tilde{S}^{03}_L, \tilde{S}^{12}_L$, $\tilde{S}^{03}_R, \tilde{S}^{12}_R$ (of $SO(1,3)$ group in the $\tilde{S}^{mn}$ sector), $\tilde{\tau}^{13}$ of $SU(2)_I$, $\tilde{\tau}^{23}$ of the second $SU(2)_{II}$, $\tilde{\tau}^4$, $\tilde{Y}' = \tilde{\tau}^{23} - \tilde{\tau}^4 \tan\tilde{\theta}_2$, taking $\tilde{\theta}^2 = 0$, $\tilde{Y} = \tilde{\tau}^4 + \tilde{\tau}^{23}$, $\tilde{Q} = \tilde{\tau}^4 + \tilde{S}^{56}$. See also the ref. [9].

the tree level.

A. Discrete symmetries of the theory unifying spin and charges and explaining families

Let us define the discrete operators of the parity (P) and of the charge conjugation (C).

$$\begin{aligned} P &= \gamma^0 \gamma^8 I_x \,, \\ C &= \Pi_{Im\,\gamma^a} \gamma^a K \,. \end{aligned} \tag{6}$$

K means complex conjugation, while in our choice of matrix representation of the γ^a matrices $\Pi_{Im\,\gamma^a} \gamma^a = \gamma^2\gamma^5\gamma^7\gamma^9\gamma^{11}\gamma^{13}$.

One can easily check that P transforms the u^{c1}_R from the first row in Table I into the u^{c1}_L of the seventh row in the same table. The CP transforms the fermion states of table I into the corresponding states of antifermions: u^{c1}_R from the first row in table I with the spin $\frac{1}{2}$, weak chargeless and of the colour charge $((\frac{1}{2}, \frac{1}{2\sqrt{3}}))$ into a right handed antiquark $\bar{u}^{\bar{c1}}_R$, weak

i		$\vert^a\psi_i>$	$\Gamma^{(1,3)}$	S^{12}	$\Gamma^{(4)}$	τ^{13}	τ^{23}	Y	Q
		Octet, $\Gamma^{(1,7)}=-1,\ \Gamma^{(6)}=1,$ of antiquarks							
1	$\bar{u}^{\bar{c1}}_{R}$	$\stackrel{03}{[-i]}\stackrel{12}{[-]}\vert\stackrel{56}{[-]}\stackrel{78}{(+)}\vert\vert\stackrel{9\,10}{[-]}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	1	$-\frac{1}{2}$	-1	$-\frac{1}{2}$	0	$-\frac{1}{6}$	$-\frac{2}{3}$
2	$\bar{u}^{\bar{c1}}_{R}$	$\stackrel{03}{(+i)}\stackrel{12}{(+)}\vert\stackrel{56}{[-]}\stackrel{78}{(+)}\vert\vert\stackrel{9\,10}{[-]}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	1	$\frac{1}{2}$	-1	$-\frac{1}{2}$	0	$-\frac{1}{6}$	$-\frac{2}{3}$
3	$\bar{d}^{\bar{c1}}_{R}$	$\stackrel{03}{[-i]}\stackrel{12}{[-]}\vert\stackrel{56}{(+)}\stackrel{78}{[-]}\vert\vert\stackrel{9\,10}{[-]}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	1	$-\frac{1}{2}$	-1	$\frac{1}{2}$	0	$-\frac{1}{6}$	$\frac{1}{3}$
4	$\bar{d}^{\bar{c1}}_{R}$	$\stackrel{03}{(+i)}\stackrel{12}{(+)}\vert\stackrel{56}{(+)}\stackrel{78}{[-]}\vert\vert\stackrel{9\,10}{[-]}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	1	$\frac{1}{2}$	-1	$\frac{1}{2}$	0	$-\frac{1}{6}$	$\frac{1}{3}$
5	$\bar{d}^{\bar{c1}}_{L}$	$\stackrel{03}{(+i)}\stackrel{12}{[-]}\vert\stackrel{56}{(+)}\stackrel{78}{(+)}\vert\vert\stackrel{9\,10}{[-]}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	-1	$-\frac{1}{2}$	1	0	$\frac{1}{2}$	$\frac{1}{3}$	$\frac{1}{3}$
6	$\bar{d}^{\bar{c1}}_{L}$	$\stackrel{03}{[-i])}\stackrel{12}{(+)}\vert\stackrel{56}{(+)}\stackrel{78}{(+)}\vert\vert\stackrel{9\,10}{[-]}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	-1	$\frac{1}{2}$	1	0	$\frac{1}{2}$	$\frac{1}{3}$	$\frac{1}{3}$
7	$\bar{u}^{\bar{c1}}_{L}$	$\stackrel{03}{(+i)}\stackrel{12}{[-]}\vert\stackrel{56}{[-]}\stackrel{78}{[-]}\vert\vert\stackrel{9\,10}{[-]}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	-1	$-\frac{1}{2}$	1	0	$-\frac{1}{2}$	$-\frac{2}{3}$	$-\frac{2}{3}$
8	$\bar{u}^{\bar{c1}}_{L}$	$\stackrel{03}{[-i]}\stackrel{12}{(+)}\vert\stackrel{56}{[-]}\stackrel{78}{[-]}\vert\vert\stackrel{9\,10}{[-]}\ \stackrel{11\,12}{(+)}\ \stackrel{13\,14}{(+)}$	-1	$\frac{1}{2}$	1	0	$-\frac{1}{2}$	$-\frac{2}{3}$	$-\frac{2}{3}$

TABLE V: The 8-plet of antiquarks to the quarks obtained from Table I by the CP ($=\gamma^2\gamma^5\gamma^7\gamma^9\gamma^{11}\gamma^{13}K\gamma^0\gamma^8\, I_x$) conjugation.

charged and of the colour charge $((-\frac{1}{2}, -\frac{1}{2\sqrt{3}}))$ as presented in table V.

III. THE FERMION-ANTIFERMION ASYMMETRY WITHIN THE THEORY UNIFYING SPIN AND CHARGES AND EXPLAINING FAMILIES

As said in the abstract, I shall here follow the ideas from the references [3–6]. The difference from the studies there in here is, as explained, in the number of families (there are two decoupled groups of four families and consequently two stable families), in the number of gauge fields contributing to the phase transitions and in the types of the gauge fields contributing to phase transitions.

Let us assume that the fermion-antifermion asymmetry is zero, when the expanding universe cools down to the temperature below the unification scale of the spin and the charges that is to the temperature below, let say, 10^{16} TeV, when there are eight massless families, manifesting the symmetry $SO(1,3) \times SU(2) \times SU(2) \times U(1) \times SU(3)$, and distinguishing among themselves in the quantum numbers defined by $\tilde{S}^{ab}$.

Then we must investigate, how much do the following processes contribute to the fermion-antifermion asymmetry in non equilibrium thermal processes in the expanding universe:

- The nonconservation of currents on the quantum level due to the triangle anomalies [4–6], which are responsible for P and CP nonconservation

$$\partial^m \, j^{Ai\,\alpha(i)}_m \;=\; \frac{(g^A)^2}{8\pi^2}\,\frac{1}{2}\,\varepsilon_{mnpr}\,F^{Ai\,mn}F^{Ai\,pr}. \tag{7}$$

Here $j^{Ai\,\alpha}_m$ stays for the currents of fermions (and antifermions), which carry a particular charge denoted by a charge group A, in our case $A = 4$ means the $U(1)$ charge originating in $SO(6)$, $A = 3$ means the $SU(3)$ (colour) charge, $A = 2_I$ means the weak $SU(2)_I$ charge of the left handed doublets, while $A = 2_{II}$ stays for the $SU(2)_{II}$ charge of the right handed singlets before the $SU(2)_{II}$ break, $A = 1$ stays for the actual $U(1)$ charge (the *standard model* like hyper charge after the $SU(2)_{II}$ break and the electromagnetic one after the weak break).

In my case also the fields, which look like scalar fields in $d = (1+3)$, $\tilde{A}^{\tilde{A}i}_s$, $s,t \in 5,6,\cdots$, and to which the fermions are coupled, contribute.

All the fermions and antifermions, which are coupled to a particular gauge field A^{Ai}_m and in my case also $\tilde{A}^{\tilde{A}i}_s$ contribute to the current

$$j^{Ai\,\alpha(i)}_m = \psi^{Ai\,\alpha(i)\dagger}\,\gamma^0\gamma^m\,\psi^{Ai\,\alpha(i)}. \tag{8}$$

$(i) \in \{1,8\}$ enumerates families, in my case twice four families which are distinguishable by the quantum numbers originating in $\tilde{S}^{ab}$, namely, after the break of $SU(2)_I$ the lower four families, which are doublets with respect to $\tilde{N}^i_+$ and $\tilde{\tau}^{I\,i}$ and singlets with respect to $\tilde{N}^i_-$ and $\tilde{\tau}^{II\,i}$, stay massless, while the upper four families are doublets with respect to $\tilde{N}^i_-$ and $\tilde{\tau}^{IIi}$ and singlets with respect to $\tilde{N}^i_+$ and $\tilde{\tau}^{Ii}$. After the electroweak break all the eight families become massive, but the upper four families have no mixing matrix elements since the way of breaking leaves all the ω_{msa} and $\tilde{\omega}_{msa}$, with $m = 0,1,2,3; s = 5,6,\cdots$, equal to zero. α distinguishes the multiplets in each family, in my case of the two $SU(2)$ gauge groups α distinguishes the $SU(2)_I$ doublets, that is one colour singlet and one colour triplet, and the $SU(2)_{II}$ doublets, again one colour singlet and one colour triplet. A^{Ai}_m are the corresponding gauge fields, with tensors $F^A_{mn} = \tau^{Ai}\,F^{Ai}_{mn}$ and $F^{Ai}_{mn} = A^{Ai}_{n,m} - A^{Ai}_{m,n} + g^A\,f^{Aijk}\,A^{Aj}_m\,A^{Ak}_n$. (The scalar fields $\tilde{A}^{\tilde{A}i}_s$ define tensors $\tilde{F}^{\tilde{A}i\,st} = \tilde{A}^{\tilde{A}i}_{t,s} - \tilde{A}^{\tilde{A}i}_{s,t} + \tilde{g}^A\,f^{A\,ijk}\tilde{A}^{\tilde{A}j}_s\,\tilde{A}^{\tilde{A}k}_t$.)

The nonconserved currents affect the fermions and antifermions. (In the later case the τ^{Ai} are replaced by $\bar{\tau}^{Ai}$, both fulfilling the same commutation relations $\{\tau^{Ai}, \tau^{Bj}\}_- = i\delta^{AB} f^{A\,ijk}\,\tau^{Ak}$, $\{\bar{\tau}^{Ai}, \bar{\tau}^{Bj}\}_- = i\delta^{AB} f^{A\,ijk}\,\bar{\tau}^{Ak}$). One obtains $\bar{\tau}^{Ai}$ from τ^{Ai} by the CP transformation $P = \gamma^0\gamma^8\, I_x$, while $C = \prod_{Im\gamma^a} \gamma^a\, K$ (See II A).

- The nonconservation of the fermion numbers originating in the complex phases of the mixing matrices of the two times 4×4 mass matrices for each member of a family, after the two successive breaks causes two phase transitions when the symmetry $SU(2) \times SU(2) \times U(1)$ breaks first to $SU(2) \times U(1)$ and finally to $U(1)$ and the two types of gauge fields manifest their masses while the two groups of four with the mixing matrices decoupled families gain nonzero mass matrices in the first break the upper four families and in the second break the lower four families.

I am following here the references [3–6]. The nonconservation of currents may be expected whenever the non-Abelian gauge fields manifest a non trivial structure of vacua, originating in the *instanton* solutions of the Euclidean non-Abelian gauge theories in $(1{+}3)-$dimensional space, that is in A^A_m, which fulfil the boundary condition $\lim_{r\to\infty}\ \tau^{Ai}\, A^{Ai}_m = U^{-1}\partial_m U$, summed over i for a particular gauge group A (and similarly might be that the fields $\lim_{\rho\to\infty}\ \tilde{\tau}^{\tilde{A}i}\, \tilde{A}^{\tilde{A}i}_s = U^{-1}\partial_s U$, with $r = \sqrt{(x^0)^2 + \vec{x}^2}$ and $\rho = \sqrt{\sum_s (x^s)^2}$, for a particular $\tilde{A}$, contribute as well, where the effect of the triangle anomalies in the case of scalar gauge fields depending on $x^\sigma, \sigma = 5, 6, 7, 8$ and the corresponding meaning of the winding numbers distinguishing among the different vacua in this case might be negligible and should be studied). The vacua distinguish among themselves in the topological quantum numbers n_A ($n_{\tilde{A}}$), determined by a particular choice of U

$$n_A = \frac{(g^A)^2}{16\pi^2} \int d^4x\, \varepsilon_{mnpr}\, Tr(F^{Amn} F^{Apr}) = \frac{(g^A)^2}{32\pi^2} \int d^4x \partial_m K^{A\,m}, \tag{9}$$

where $K^A_m = \sum_i 4\varepsilon_{mnpr}\, (A^{Ai}_n \partial_p A^A_r + \frac{2}{3}\, g^A f^{Aijk} A^{Ai}_n A^{Aj}_p A^{Ak}_r)$. (Similarly also the topological quantum number $n_{\tilde{A}}$ might be non negligible.)

Instanton solutions fulfilling the boundary condition $\lim_{r\to\infty}\ \tau^{Ai}\, A^{Ai}_m = U^{-1}\partial_m U$ for a particular gauge group A (or $\lim_{\rho\to\infty}\ \tilde{\tau}^{\tilde{A}i}\, \tilde{A}^{\tilde{A}i}_s = U^{-1}\partial_m U$ for a particular $\tilde{A}$), each with its own U for a particular A (or $\tilde{A}$), connect vacua $|n_A>$ with different winding numbers [17] n_A (and correspondingly for $n_{\tilde{A}}$). The true vacuum $|\theta^A>$ is for each A (let it count also $\tilde{A}$) in a stationary situation a superposition of the vacua, determined by the time independent gauge transformation [3] $\mathcal{T}$, $\mathcal{T}|\theta^A >= e^{i\theta^A}|\theta^A >$, where θ^A is a parameter, which weights the contribution of a vacuum to the effective Lagrange density $\mathcal{L}_{eff} = \mathcal{L} + \sum_A \frac{\theta^A}{16\pi^2} F^{Ai\,mn} \frac{1}{2}\varepsilon_{mnpr} F^{Ai\,pr}$, for a particular gauge field. $\mathcal{T}$ acts as the raising operator for the handedness (chirality). The second term of the effective Lagrange density $\mathcal{L}_{eff}$ violates parity P and then also CP. The vacuum state with the definite handedness has also a definite topological quantum number. In the presence of the massless fermions all the vacua $|\theta^A >$, for each A, are equivalent.

The fermion currents (Eq.(8)) are not conserved in processes, for which the gauge fields are such that the corresponding winding number n_A of Eq. (9) is nonzero. Correspondingly also the fermion (and antifermion numbers), carrying the corresponding charge, are not conserved

$$\Delta n_{Ai\,\alpha(i)} = n_A. \tag{10}$$

The fermion number of all the fermions interacting with the same non-Abelian gauge field with nonzero winding number, either of a vector or of a scalar type (whose contribution should be studied and hopefully understood), changes in such processes for the *same amount*: Any member of a family, interacting with the particular field and therefore also the corresponding members of each family, either a quark or a lepton member of doublets, change for the same amount, before the breaks or after the breaks (in my case first from $SO(1,3) \times SU(2) \times SU(2) \times U(1) \times SU(3)$ to $SO(1,3) \times SU(2) \times U(1) \times SU(3)$ and finally to $SO(1,3) \times U(1) \times SU(3)$) of the symmetries.

For a baryon three quarks are needed. It is the conservation of the colour charge which requires that the lepton number and the baryon number ought to be conserved separately as long as the charge group is a global symmetry. The transformations, which allow rotations of a lepton to a quark or opposite, conserve the fermion number, but not the lepton and not the baryon number.

Instanton solutions of the non-Abelian gauge fields, which connect different vacua (see the refs. [6], page 481, and [4], page 6), are characterized by the highest value of the *instanton* field between the two vacua, that is by the *sphaleron* energy.

The question arises, can the *instanton* solutions be responsible for the baryon asymmetry of the universe? The authors of the papers [4, 5] discuss and evaluate the probability for tunnelling from one vacuum to the other at low energy regime and also at the energies of *sphalerons*. When once the system of gauge fields is in one vacuum the probability for the transition to another vacuum depends not only on the *sphalerons* height (energy) but also on the temperature. If the temperature is low, then the transition is negligible. At the temperature above the phase transition (the authors [4] discuss the electroweak phase transition starting at around 100 GeV, while in my case there is also the $SU(2)_{II}$ phase transition at around 10^{16} GeV or slightly below) when the fermions are massless and the expansion rate of the universe is much slower that the rate of nonconservation of the fermion number, and in the case of non equilibrium processes in phase transitions, the fermion number nonconservation can be large. The authors conclude that more precise evaluations (treating several models) of the probability that in a non thermal equilibrium phase transition and below it the fermion number would not be conserved due to transitions to vacua with different winding numbers in the amount as observed for the (first family) baryon number excess in the universe are needed.

What can be concluded about the fermion number asymmetry, caused by the transitions of gauge fields to different vacua, in my case, where at energies above the $SU2_{II}$ phase transition there are eight families of massless fermions, with the charges manifesting the symmetries first of $SU(2)_I \times SU(2)_{II} \times U(1)$ and correspondingly with the two kinds of the vector gauge $SU(2)$ fields which both might demonstrate the vacua with different winding numbers? In addition also the scalar gauge fields might contribute with their even more rich vacua (if they do that at all). The phase transitions caused first by the break of the symmetry $SU(2)_I \times SU(2)_{II} \times U(1)$ to $SU(2)_{II} \times U(1)$, when the upper four families gain masses (and the corresponding gauge vector fields become massive) and then by the final break to $U(1)$, with the $\tilde{S}^{ab}$ sector causing the masses in both transitions and may be also taking care of the richness of vacua with different winding numbers, might show up after a careful study as a mechanism for generating the fermion-antifermion (or the antifermion-fermion) asymmetry. Although I do not yet see, how do the non equilibrium processes in the first order phase transitions decide about the excess of either fermions or of antifermions.

So, is it in my case possible that the two successive non equilibrium phase transitions leave the excess of antifermions in the case of the upper four families and the excess of fermions in the lower four families? Or there is a negligible excess of either fermions or antifermions in the upper four families? We saw in the ref. [10] that an excess of either fermions or antifermions is not important for massive enough (few 100 TeV) stable fifth family members. The excess of fermions over antifermions is certainly what universe made a choice of for the lower four families, whatever the reason for this fact is. Can this be easier understood within the *spin-charge-family-theory*? All these need a careful study.

The fermion number nonconservation originates also in the complex phases of the mass mixing matrices of each of the two groups of four family members. It might be that the vacua, triggered by *instanton* solutions of the gauge vector and scalar fields, and the mass matrices, determined on the tree level by the vacuum expectation values of the scalar gauge fields in the $\tilde{S}^{ab}$ sector, are connected (since in the *instanton* case also the scalar fields, the gauge fields of charges originating in $\tilde{S}^{ab}$ might exhibit the *instanton* solutions).

IV. CONCLUSION

In this contribution I pay attention to the origin of baryon asymmetry of our universe within the *spin-charge-family-theory* under the assumption that the asymmetry is caused **i.** by the *instanton* solutions of the non-Abelian gauge fields which determine vacua with different winding numbers and **ii.** by the complex matrix elements of the mixing matrices.

The *spin-charge-family-theory* namely assumes besides the Dirac Clifford algebra objects also the second ones $\tilde{\gamma}^a$ as a necessary mechanism (or better a mathematical tool) which should be used in order that we consistently describe both: spin and charges, as well as families. The second kind is namely responsible for generating families, defining the equivalent representations with respect to the Dirac one. Correspondingly there are besides the two kinds of the vector gauge fields, the $SU(2)_I$ and $SU(2)_{II}$, also the scalar gauge fields, the two $SU(2)$ from $SO(4)$ and the two $SU(2)$ from $SO(1,3)$, the superposition of the gauge fields of $\tilde{S}^{ab}(=\frac{i}{4}(\tilde{\gamma}^a\tilde{\gamma}^b-\tilde{\gamma}^b\tilde{\gamma}^a))$, which might contribute to vacua with different winding numbers (what has to be studied). The scalar fields, originating in the $\tilde{S}^{ab}$ charges, are responsible with their vacuum expectation values (and in loop corrections) for the mass matrices of fermions after the breaks of symmetries.

The theory predicts twice four families (which differ in the family quantum numbers in the way that the upper four families are doublets with respect to $\tilde{\tau}^{II\,i}$ and $\tilde{N}^i_-$, while the lower four families are doublets with respect to $\tilde{\tau}^{I\,i}$, and $\tilde{N}^i_+$) which all are massless above the last two phase transitions.

What should be clarified in the *spin-charge-family-theory* is whether the predicted twice four families (rather than once three families of the *standard model*) and the fact that there are gauge fields belonging to two kinds of generator (S^{ab} and $\tilde{S}^{ab}$) make the baryon number asymmetry easier to be understood within these two phenomena — the *instanton* responsibility for the fermion number nonconservation and the complex matrix elements of the mixing matrices responsibility for the fermion number nonconservation.

The manifestation of the *instanton* gauge vector and scalar fields in the determination of the properties of the vacuum might be correlated with the vacuum expectation values of the scalar fields defining the mass matrices of twice the four families. Both manifestations appear in possibly non equilibrium phase transitions of the expanding universe, which cause breaking of particular symmetries and also the fermion number nonconservation. In this contribution I just follow the way suggested by the ref. [4] and by the authors cited in this reference, while taking into account the requirement of the *spin-charge-family-theory*. The fermion number nonconservation obviously ended in the excess of (what we call) fermions for the lower four families, while for the upper four families we have to see whether there is the excess of either the stable fifth family fermions or antifermions. To answer these questions a careful study is needed. It even might be that there was at the non equilibrium phase transitions the same excess of antifermions for the upper four families as it is of fermions for the lower four families, while later the complex matrix elements in the mixing matrices change this equality drastically. But yet it must be understood the origin of both sources of the fermion number nonconservation.

Appendix: Some useful relations

The following Cartan subalgebra set of the algebra S^{ab} (for both sectors) is chosen:

$$\begin{aligned}
&S^{03}, S^{12}, S^{56}, S^{78}, S^{9\,10}, S^{11\,12}, S^{13\,14}\\
&\tilde{S}^{03}, \tilde{S}^{12}, \tilde{S}^{56}, \tilde{S}^{78}, \tilde{S}^{9\,10}, \tilde{S}^{11\,12}, \tilde{S}^{13\,14}.
\end{aligned} \tag{A.1}$$

A left handed ($\Gamma^{(1,13)} = -1$) eigen state of all the members of the Cartan subalgebra

$$\begin{aligned}
&\stackrel{03}{(+i)}\stackrel{12}{(+)} \mid \stackrel{56}{(+)}\stackrel{78}{(+)} \mid\mid \stackrel{9\,10}{(+)}\stackrel{11\,12}{(-)}\stackrel{13\,14}{(-)} \, |\psi\rangle = \\
&\frac{1}{2^7}(\gamma^0 - \gamma^3)(\gamma^1 + i\gamma^2)|(\gamma^5 + i\gamma^6)(\gamma^7 + i\gamma^8)|| \\
&(\gamma^9 + i\gamma^{10})(\gamma^{11} - i\gamma^{12})(\gamma^{13} - i\gamma^{14})|\psi\rangle.
\end{aligned} \tag{A.2}$$

represent the u_R-quark with spin up and of one colour.

$\tilde{S}^{ab}$ generate families from the starting u_R quark In particular $\tilde{S}^{03}(=\frac{i}{2}[\stackrel{03}{(\tilde{+i})}\stackrel{12}{(\tilde{+})} + \stackrel{03}{(\tilde{-i})}\stackrel{12}{(\tilde{+})} + \stackrel{03}{(\tilde{+i})}\stackrel{12}{(\tilde{-})} + \stackrel{03}{(\tilde{-i})}\stackrel{12}{(\tilde{-})}])$ applied on a right handed u_R-quark with spin up and a particular colour generate a state which is again a right handed u-quark of the same colour.

$$
\begin{aligned}
\stackrel{03}{(\tilde{-i})}\stackrel{12}{(\tilde{-})} \quad & \stackrel{03}{(+i)}\stackrel{12}{(+)} \mid \stackrel{56}{(+)}\stackrel{78}{(+)} \mid\mid \stackrel{910}{(+)}\stackrel{1112}{(-)}\stackrel{1314}{(-)} = \\
& \stackrel{03}{[+i]}\stackrel{12}{[+]} \mid \stackrel{56}{(+)}\stackrel{78}{(+)} \mid\mid \stackrel{910}{(+)}\stackrel{1112}{(-)}\stackrel{1314}{(-)},
\end{aligned} \tag{A.3}
$$

where

$$
\begin{aligned}
\stackrel{ab}{(\pm i)} &= \frac{1}{2}\,(\gamma^a \mp \gamma^b), \stackrel{ab}{(\pm 1)} = \frac{1}{2}\,(\gamma^a \pm i\gamma^b), \\
\stackrel{ab}{[\pm i]} &= \frac{1}{2}(1 \pm \gamma^a\gamma^b), \quad \stackrel{ab}{[\pm 1]} = \frac{1}{2}(1 \pm i\gamma^a\gamma^b), \\
\stackrel{ab}{(\tilde{\pm i})} &= \frac{1}{2}(\tilde{\gamma}^a \mp \tilde{\gamma}^b), \quad \stackrel{ab}{(\tilde{\pm 1})} = \frac{1}{2}(\tilde{\gamma}^a \pm i\tilde{\gamma}^b), \\
\stackrel{ab}{[\tilde{\pm i}]} &= \frac{1}{2}(1 \pm \tilde{\gamma}^a\tilde{\gamma}^b), \quad \stackrel{ab}{[\tilde{\pm 1}]} = \frac{1}{2}(1 \pm i\tilde{\gamma}^a\tilde{\gamma}^b).
\end{aligned} \tag{A.4}
$$

We present below some useful relations which are easy to derive [2].

$$
\begin{aligned}
\stackrel{ab}{(k)}\stackrel{ab}{(k)} &= 0, \; \stackrel{ab}{(k)}\stackrel{ab}{(-k)} = \eta^{aa}\;\stackrel{ab}{[k]}, \; \stackrel{ab}{[k]}\stackrel{ab}{[k]} = \stackrel{ab}{[k]}, \\
\stackrel{ab}{[k]}\stackrel{ab}{[-k]} &= 0, \stackrel{ab}{(k)}\stackrel{ab}{[k]} = 0, \quad \stackrel{ab}{[k]}\stackrel{ab}{(k)} = \stackrel{ab}{(k)}, \\
\stackrel{ab}{(k)}\stackrel{ab}{[-k]} &= \stackrel{ab}{(k)}, \quad \stackrel{ab}{[k]}\stackrel{ab}{(-k)} = 0.
\end{aligned} \tag{A.5}
$$

$$
\begin{aligned}
\stackrel{ab}{(\tilde{k})}\stackrel{ab}{(k)} &= 0, \quad \stackrel{ab}{(\tilde{-k})}\stackrel{ab}{(k)} = -i\eta^{aa}\;\stackrel{ab}{[k]}, \\
\stackrel{ab}{(\tilde{k})}\stackrel{ab}{[k]} &= i\;\stackrel{ab}{(k)}, \quad \stackrel{ab}{(\tilde{k})}\stackrel{ab}{[-k]} = 0.
\end{aligned} \tag{A.6}
$$

$$
\begin{aligned}
N^{\pm}_{+} &= N^1_{+} \pm i\,N^2_{+} = -\stackrel{03}{(\mp i)}\stackrel{12}{(\pm)}, \quad N^{\pm}_{-} = N^1_{-} \pm i\,N^2_{-} = \stackrel{03}{(\pm i)}\stackrel{12}{(\pm)}, \\
\tilde{N}^{\pm}_{+} &= -\stackrel{03}{(\tilde{\mp i})}\stackrel{12}{(\tilde{\pm})}, \quad \tilde{N}^{\pm}_{-} = \stackrel{03}{(\tilde{\pm i})}\stackrel{12}{(\tilde{\pm})}, \\
\tau^{1\pm} &= (\mp)\;\stackrel{56}{(\pm)}\stackrel{78}{(\mp)}, \quad \tau^{2\pm} = (\mp)\;\stackrel{56}{(\mp)}\stackrel{78}{(\mp)}, \\
\tilde{\tau}^{1\pm} &= (\mp)\;\stackrel{56}{(\tilde{\pm})}\stackrel{78}{(\tilde{\mp})}, \quad \tilde{\tau}^{2\pm} = (\mp)\;\stackrel{56}{(\tilde{\mp})}\stackrel{78}{(\tilde{\mp})}\;.
\end{aligned} \tag{A.7}
$$

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$(x^0)^2 + \vec{x}^2$, is presented. Operators $\vec{\sigma}^A$ are the three Pauli matrices, used to denote the $SU(2)$ gauge group in this case: $\vec{\tau}^A = \frac{\vec{\sigma}^A}{2}$, $\{\tau^{Ai}, \tau^{Aj}\}_- = i\varepsilon^{ijk}\tau^{Ak}$. The corresponding action $\int d^4x \, \frac{1}{2}\varepsilon_{mnpr} F^{Ai\,mn} F^{Aipr} = \frac{8\pi^2}{g^2}$, while $U = \frac{x^0 + i\vec{\sigma}^A \cdot \vec{x}}{r}$, defines the $n = 1$ vacuum state.