The *approach unifying spins and charges* is offering a new way beyond the Standard model: A simple action, in which in $d > 1 + 3$ spinors carry only two kinds of a spin, no charges, manifests in $d = 1 + 3$ the Standard model effective Lagrangean—with the known gauge fields, families and their mass matrices—predicting the fourth family and the stable fifth family as the Dark matter candidate.

N.S. Mankoč Borštnik, Bled 2010, July 12-22, Ljubljana

Collaborators on this project, which **S**usana **N**orma **M**ankoč **B**orštnik has started almost 15 years ago:
Anamarija Pika **Borštnik Bračič**, Gregor **Bregar**, Matjaž **Breskvar**, Dragan **Lukman**, Holger Bech **Nielsen** (first of all), others.

- *Phys. Lett.* **B 292**, 25-29 (1992),
- *Modern Phys. Lett.* **A 10**, 587-595 (1995),
- *Int. J. Theor. Phys.* **40**, 315-337 (2001).
  Phys. Rev. **D 62** (04010-14) (2000),
- *Phys. Lett.* **B 633** (2006) 771-775, **B 644** (2007) 198-202, **B** (2008) 110.1016, (2006), hep-th/0311037, hep-th/0509101, with H.B.N.
- hep-ph/0401043, hep-ph/0401055, hep-ph/0301029,
- *Phys. Rev.* , **D 74** 073013-16 (2006), hep-ph/0512062, with A.B.B..
  hep-ph/0606159, with M.B., D.L..
- *New Jour. of Phys.* **10** (2008) 093002, hep-ph/0606159, hep/ph-07082846, with G.B., M.B., D.L.
- *Phys. Rev.* **D** (2009) 80.083534, astro-ph¿arXiv: 0907.0196, with G.B.

WHAT DOES the APPROACH OFFER and HOW DOES IT?

The **Approach unifying spins and charges** is offering **a new way beyond the Standard model** of the electroweak and colour interactions assuming that:

- A spinor carries only **two kinds of the spin, no charges** and interacts in $d = 1 + 13$ with the **vielbeins** and the two kinds of the **spin connections**.

- The **Dirac spin** **takes care in** $d = 1 + 3$ **of** **the spin and all the charges** of quarks and leptons, the **second kind of spin generates families**.
- **A simple action in** $d = 1 + 13$ **manifests** in $d = 1 + 3$, after appropriate breaks of the starting symmetry:
  the mass matrices of quarks and leptons, and the **observed gauge fields**.

We are looking for

- **general proofs that this Approach does lead in the observable (low) energy region to the observable phenomena**,
- **predictions of the Approach**.

We are trying to understand:

1. Where do families of quarks and leptons come from? Why have we observed so far three families?
2. What does determine the strength of the Yukawa couplings and the weak scale, that is where do the mass matrices originate?
3. What does determine the colour scale?
4. Why does the weak charge depend on the handedness? What does form the scalar (Higgs) field, where do masses of the massive bosons originate?
5. What does constitute the dark matter? If a new stable family, how does such a stable family behave in the evolution of the universe?
6. And others.

The **Approach unifying spins and charges might have a chance to offer the answers to these questions**:

- The representation of one Weyl spinor of the group SO(1,13), **manifests the left handed weak charged quarks and leptons and the right handed weak chargeless quarks and leptons**.
- There are (only) two kinds of the Clifford algebra objects. **One kind takes care of the spin and the charges**, **the other of the families**.
- A simple starting Lagrange density in $d = (1 + 13)$ for gauge fields (linear in the two curvatures) and a spinor (which carries two kinds of the spins and interacts with the gravitational fields—**the vielbeins** and the **two kinds** of **spin connections**—manifests in $d = (1 + 3)$ the known effective Lagrange density for the **families of spinors** and the **known**

- It is a part of a simple starting Lagrange density for a spinor in $d = (1 + 13)$, which **manifests in $d = (1 + 3)$ the mass matrices**.
- **The way of breaking symmetries determines the charges and the properties of families, as well as the coupling constants of the gauge fields**.
- There are **two times four families** with zero Yukawa couplings among the members of the first and the second group of families. The three from the lowest four families are the observed ones, the **fourth family might** (as the first rough estimations show) **be seen at the LHC**. The lowest among the **decoupled** four families is the **candidate** for forming the **Dark matter** clusters.

The ACTION of the APPROACH

There are **two kinds of the Clifford algebra objects which determine the properties of spinors (fermions)**:

- The **Dirac γ^a operators** (used by Dirac 80 years ago),
- The **second one:** $\tilde{\gamma}^a$, which I recognized in Grassmann space

$$
\begin{aligned}
\{\gamma^{\mathbf{a}}, \gamma^{\mathbf{b}}\}_+ &= \mathbf{2}\eta^{\mathbf{ab}} = \{\tilde{\gamma}^{\mathbf{a}}, \tilde{\gamma}^{\mathbf{b}}\}_+,\\
\{\gamma^{\mathbf{a}}, \tilde{\gamma}^{\mathbf{b}}\}_+ &= 0,\\
(\tilde{\gamma}^{\mathbf{a}}\mathbf{B} : &= \mathbf{i}(-)^{\mathbf{n_B}}\mathbf{B}\gamma^{\mathbf{a}}\,)\,|\psi_0>,\\
\mathbf{B} &= a_0 + a_a\gamma^a + a_{ab}\gamma^a\gamma^b + \cdots + a_{a_1\cdots a_d}\gamma^{a_1}\ldots\gamma^{a_d}
\end{aligned}
$$

$(-)^{n_B} = +1$ or -1, when the object B has a Clifford even or odd character, respectively.
Both are used in the Approach to determine properties of spinors.

$$\mathbf{S}^{ab} := (i/4)(\gamma^a\gamma^b - \gamma^b\gamma^a),$$

$$\tilde{\mathbf{S}}^{ab} := (i/4)(\tilde{\gamma}^a\tilde{\gamma}^b - \tilde{\gamma}^b\tilde{\gamma}^a),$$

$$\{\mathbf{S}^{ab}, \tilde{\mathbf{S}}^{cd}\}_- = 0.$$

$\tilde{\mathbf{S}}^{ab}$ define the equivalent representations with respect to $\mathbf{S}^{ab}$.

- I recognized: If γ^a **describe the spins and the charges of spinors,**
  describe $\tilde{\gamma}^a$ **their families**.

A simple action for **a spinor which carries in** $d = (1+13)$ **only two kinds of a spin** (no charges) and for **the gauge fields**

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

The Einstein action for a free gravitational field is assumed to be linear in the curvature

$$\begin{aligned}\mathcal{L}_g &= E\,(\alpha_\omega\,R + \tilde{\alpha}\tilde{R}),\\ R &= f^{\alpha[a}f^{\beta b]}\,(\omega_{ab\alpha,\beta} - \omega_{ca\alpha}\omega^{c}{}_{b\beta}),\\ \tilde{R} &= f^{\alpha[a}f^{\beta b]}\,(\tilde{\omega}_{ab\alpha,\beta} - \tilde{\omega}_{ca\alpha}\tilde{\omega}^{c}{}_{b\beta}),\end{aligned}$$

with $E = \det(e^{a}{}_{\alpha})$
and $f^{\alpha[a}f^{\beta b]} = f^{\alpha a}f^{\beta b} - f^{\alpha b}f^{\beta a}$.

Variation of the action brings for $\omega_{ab\alpha}$

$$
\begin{aligned}
\omega_{ab\alpha} \;=\; & -\frac{1}{2E}\Big\{ e_{e\alpha} e_{b\gamma}\, \partial_\beta (E f^{\gamma[e} f^{\beta}{}_{a]}) + e_{e\alpha} e_{a\gamma}\, \partial_\beta (E f^{\gamma}{}_{[b} f^{\beta e]}) \\
& \qquad - e_{e\alpha} e^{e}{}_{\gamma}\, \partial_\beta \big(E f^{\gamma}{}_{[a} f^{\beta}{}_{b]}\big) \Big\} \\
& - \frac{e_{e\alpha}}{4} \Big\{ \bar{\Psi} \Big(\gamma_e S_{ab} + \frac{3i}{2} (\delta^{e}_{b} \gamma_a - \delta^{e}_{a} \gamma_b) \Big) \Psi \Big\} \\
& - \frac{1}{d-2} \Big\{ e_{a\alpha} \Big[\frac{1}{E} e^{d}{}_{\gamma} \partial_\beta \Big(E f^{\gamma}{}_{[d} f^{\beta}{}_{b]} \Big) + \frac{1}{2} \bar{\Psi} \gamma^{d} S_{db} \Psi \Big] \\
& \qquad - e_{b\alpha} \Big[\frac{1}{E} e^{d}{}_{\gamma} \partial_\beta \Big(E f^{\gamma}{}_{[d} f^{\beta}{}_{a]} \Big) + \frac{1}{2} \bar{\Psi} \gamma^{d} S_{da} \Psi \Big\} \Big]
\end{aligned}
$$

and for $\tilde{\omega}_{ab\alpha}$

$$
\begin{aligned}
\tilde{\omega}_{ab\alpha} &= -\frac{1}{2E}\Big\{ e_{e\alpha} e_{b\gamma}\, \partial_\beta (E f^{\gamma[e} f^{\beta}{}_{a]}) + e_{e\alpha} e_{a\gamma}\, \partial_\beta (E f^{\gamma}{}_{[b} f^{\beta e]}) \\
&\qquad - e_{e\alpha} e^{e}{}_{\gamma}\, \partial_\beta \big(E f^{\gamma}{}_{[a} f^{\beta}{}_{b]}\big) \Big\} \\
&- \frac{e_{e\alpha}}{4} \Big\{ \bar{\Psi} \Big(\gamma_e \tilde{S}_{ab} + \frac{3i}{2} (\delta^e_b \gamma_a - \delta^e_a \gamma_b) \Big) \Psi \Big\} \\
&- \frac{1}{d-2} \Big\{ e_{a\alpha} \Big[\frac{1}{E} e^{d}{}_{\gamma} \partial_\beta \Big(E f^{\gamma}{}_{[d} f^{\beta}{}_{b]} \Big) + \frac{1}{2} \bar{\Psi} \gamma^d\, \tilde{S}_{db} \Psi \Big] \\
&\qquad - e_{b\alpha} \Big[\frac{1}{E} e^{d}{}_{\gamma} \partial_\beta \Big(E f^{\gamma}{}_{[d} f^{\beta}{}_{a]} \Big) + \frac{1}{2} \bar{\Psi} \gamma^d\, \tilde{S}_{da} \Psi \Big] \Big\}
\end{aligned}
$$

The action for spinors can formally be rewritten as

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

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

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

If there is no gravity in $d = (1+3)$ the vielbeins together with the two kinds of the spin connection fields are expected to manifest, after the break of symmetries, the effective action

$$
\begin{aligned}
S_b \; = \; \int \; d^{(1+3)}x \, \{ & -\frac{\varepsilon^A}{4} \, \mathbf{F^{Aimn}} \, \mathbf{F^{Ai}}_{\mathbf{mn}} \\
& + \frac{1}{2} \, (m^{Ai})^2 \; \mathbf{A^{Ai}_m} \, \mathbf{A^{Ai\,m}} \\
& + \text{scalar terms} \, \}.
\end{aligned}
$$

The summation over A, i is assumed.

$$
\begin{array}{llll}
\mathbf{A} & = & 1 & \mathbf{U(1)}\ \text{hyper charge}\quad i=\{1\} \qquad \cdots\cdots \text{usually}\ \mathbf{Y},\\
\mathbf{A} & = & 2 & \cdots\mathbf{SU(2)}\ \text{weak charge}\quad i=\{1,2,3\} \qquad \cdots \text{usually}\ \tau^{\mathbf{i}},\\
\mathbf{A} & = & 3 & \mathbf{SU(3)}\ \text{colour charge}\quad i=\{1,\cdots,8\} \quad \cdots \text{usually}\ \lambda^{\mathbf{i}}/\mathbf{2},
\end{array}
$$

The papers, which demonstrate the possibilities, that the break of symmetries can lead to massless fermions:

- hep-th/January 2009, with H.B. Nielsen, D. Lukman,
- *Phys. Lett.* **B 644** (2007)198-202, hep-th/0608006, coauthor: H.B. Nielsen,
- *Phys. Lett.* **B** 10.1016 (2008), hep-th/0612126, arxiv:07.10.1956, with H.B. Nielsen,
- *Phys. Lett.* **B 633** (2006) 771-775, hep-th/0311037, hep-th/0509101, with H.B. Nielsen,
- arXiv:0912.4532, p. 103-118, arXiv:1001.4679v3, with H.B. Nielsen, D. Lukman.

There are **particular breaks —the particular isomertries— of the starting symmetries which make that only the measured gauge fields manifest at low energies**.

One can explicitly see that:

One Weyl spinor representation in $d = (1+13)$ with the

spin, determined by S^{ab} **as the only internal degree of freedom of one family**,

manifests, if analysed in terms of the subgroups

$SO(1,3) \times U(1) \times SU(2) \times SU(3)$,

in four-dimensional **"physical"** space

as the **ordinary** $(SO(1,3))$ **spinor** with **all the known charges** of **one family** of the **left handed weak charged** and the **right handed weak chargeless** quarks and leptons of the Standard model, with the right handed neutrinos included.

The **second kind** of the Clifford algebra objects $\tilde{S}^{ab}$

takes care of the **families**

by generating the equivalent representations with respect to S^{ab}.

THE YUKAWA COUPLINGS, SCALAR AND GAUGE FIELDS

It is a **part** of the simple **starting action for spinors in** $d = 1 + 13$ which manifests in $d = 1 + 3$ on the tree level the **mass matrices**.

It is a **part** of the simple **starting vielbeins in** $d = 1 + 13$ which manifests in $d = 1 + 3$ at low energies the **the masses of the gauge bosons**.

It is a **part** of the simple **starting vielbeins and spin connections in** $d = 1 + 13$ which manifests in $d = 1 + 3$ the **gauge fields**.

The symmetry $SO(1,7) \times U(1)$ **breaks in two steps**:

$SO(1,7) \times U(1)$ into $SO(1,3) \times SU(2) \times U(1)$

leading to the Standard model **massless quarks and leptons** of **four** (not three) families

and **massive**, **decoupled** in the Yukawa couplings from the lower mass families, **four families**.

The Yukawa couplings

$$
\begin{aligned}
-\mathcal{L}_Y &= \psi^\dagger \gamma^{\mathbf{0}} \gamma^{\mathbf{s}} \mathbf{p_{0s}} \psi \\
&= \psi^\dagger \, \gamma^{\mathbf{0}} \{ \stackrel{\mathbf{78}}{(+)} \; \mathbf{p_{0+}} + \stackrel{\mathbf{78}}{(-)} \; \mathbf{p_{0-}} \} \psi ,
\end{aligned}
$$

$$
\begin{aligned}
\mathbf{p_{0\pm}} &= (\mathbf{p_7} \mp \mathbf{i} \, \mathbf{p_8}) - \\
& \frac{1}{2} \mathbf{S^{ab}} \omega_{\mathbf{ab\pm}} - \frac{1}{2} \tilde{\mathbf{S}}^{\mathbf{ab}} \tilde{\omega}_{\mathbf{ab\pm}} ,
\end{aligned}
$$

$$
\begin{aligned}
\omega_{ab\pm} &= \omega_{ab7} \mp i \, \omega_{ab8} , \\
\tilde{\omega}_{ab\pm} &= \tilde{\omega}_{ab7} \mp i \, \tilde{\omega}_{ab8}
\end{aligned}
$$

We put $p_7 = p_8 = 0$.

The **vielbeins** together with the two kinds of **spin connections** in $d > (1+3)$ manifest in $d = (1+3)$ the **gauge fields and the masses of the gauge fields**.

$$e^{a}{}_{\alpha} = \begin{pmatrix} \delta^{m}{}_{\mu} & e^{m}{}_{\sigma} = 0 \\ e^{s}{}_{\mu} = e^{s}{}_{\sigma} \mathbf{E}^{\sigma}{}_{\mathbf{Ai}} \mathbf{A}^{\mathbf{Ai}}_{\mu} & \mathbf{e}^{\mathbf{s}}{}_{\sigma} \end{pmatrix}$$

Here $\mathbf{E}^{\sigma}{}_{\mathbf{Ai}}\, \mathbf{A}^{\mathbf{Ai}}_{\mu}$ stays for $\tau^{\mathbf{Ai}} x^{\sigma}\; \mathbf{A}^{\mathbf{Ai}}_{\mu}$.

A= 1 $\cdots$ the **U(1) field**,

A= 2 $\cdots$ the **weak field**

A= 3 $\cdots$ the **colour field**.

To the break of symmetries both kinds of the spin connection fields, together with the vielbeins, all of them manifesting as scalar fields in $d-(1+3)$, contribute.

Our technique to represent spinor states

Our technique to represent spinors works elegantly.

- *J. of Math. Phys.* **43**, 5782-5803 (2002), hep-th/0111257,
- *J. of Math. Phys.* **44** 4817-4827 (2003), hep-th/0303224, both with H.B. Nielsen.

$$\begin{aligned}
\overset{ab}{(\pm \mathbf{i})} : &= \frac{1}{2}(\gamma^{\mathbf{a}} \mp \gamma^{\mathbf{b}}),\ \overset{ab}{[\pm \mathbf{i}]} := \frac{1}{2}(1 \pm \gamma^{\mathbf{a}}\gamma^{\mathbf{b}}) \\
&\quad \text{for } \eta^{aa}\eta^{bb} = -1, \\
\overset{ab}{(\pm)} : &= \frac{1}{2}(\gamma^{\mathbf{a}} \pm \mathbf{i}\gamma^{\mathbf{b}}),\ \overset{ab}{[\pm]} := \frac{1}{2}(1 \pm i\gamma^{\mathbf{a}}\gamma^{\mathbf{b}}), \\
&\quad \text{for } \eta^{aa}\eta^{bb} = 1
\end{aligned}$$

with γ^a which are the usual **Dirac operators**

Our technique

$$\mathbf{S}^{\mathbf{ab}} \overset{\mathbf{ab}}{(\mathbf{k})} = \frac{k}{2} \overset{\mathbf{ab}}{(\mathbf{k})}, \quad \mathbf{S}^{\mathbf{ab}} \overset{\mathbf{ab}}{[\mathbf{k}]} = \frac{k}{2} \overset{\mathbf{ab}}{[\mathbf{k}]},$$

$$\tilde{\mathbf{S}}^{\mathbf{ab}} \overset{\mathbf{ab}}{(\mathbf{k})} = \frac{k}{2} \overset{\mathbf{ab}}{(\mathbf{k})}, \quad \tilde{\mathbf{S}}^{\mathbf{ab}} \overset{\mathbf{ab}}{[\mathbf{k}]} = -\frac{k}{2} \overset{\mathbf{ab}}{[\mathbf{k}]}.$$

$$\gamma^{\mathbf{a}} \overset{\mathbf{ab}}{(\mathbf{k})} = \eta^{aa} \overset{\mathbf{ab}}{[-\mathbf{k}]}, \quad \gamma^{\mathbf{b}} \overset{\mathbf{ab}}{(\mathbf{k})} = -ik \overset{\mathbf{ab}}{[-\mathbf{k}]},$$

$$\gamma^{\mathbf{a}} \overset{\mathbf{ab}}{[\mathbf{k}]} = \overset{\mathbf{ab}}{(-\mathbf{k})}, \quad \gamma^{\mathbf{b}} \overset{\mathbf{ab}}{[\mathbf{k}]} = -ik\eta^{aa} \overset{\mathbf{ab}}{(-\mathbf{k})}$$

$$\tilde{\gamma}^{\mathbf{a}} \overset{\mathbf{ab}}{(\mathbf{k})} = -i\eta^{aa} \overset{\mathbf{ab}}{[\mathbf{k}]}, \quad \tilde{\gamma}^{\mathbf{b}} \overset{\mathbf{ab}}{(\mathbf{k})} = -k \overset{\mathbf{ab}}{[\mathbf{k}]},$$

$$\tilde{\gamma}^{\mathbf{a}} \overset{\mathbf{ab}}{[\mathbf{k}]} = i \overset{\mathbf{ab}}{(\mathbf{k})}, \quad \tilde{\gamma}^{\mathbf{b}} \overset{\mathbf{ab}}{[\mathbf{k}]} = -k\eta^{aa} \overset{\mathbf{ab}}{(\mathbf{k})}.$$

Our technique

γ^a **transforms** $\overset{ab}{(k)}$ **into** $\overset{ab}{[-k]}$**, never to** $\overset{ab}{[k]}$**.**

$\tilde{\gamma}^a$ transforms $\overset{ab}{(k)}$ **into** $\overset{ab}{[k]}$**, never to** $\overset{ab}{[-k]}$**.**

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

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

$$
\stackrel{ab}{(\widetilde{\pm \mathbf{i}})} = \frac{1}{2}(\tilde{\gamma}^{\mathbf{a}} \mp \tilde{\gamma}^{\mathbf{b}}), \quad \stackrel{ab}{(\widetilde{\pm \mathbf{1}})} = \frac{1}{2}(\tilde{\gamma}^{\mathbf{a}} \pm \mathbf{i}\tilde{\gamma}^{\mathbf{b}}),
$$

The representation of a left handed Weyl spinor in $d = 1 + 13$, if analysed in terms of the "standard model" symmetries

The representations of families, scalar and gauge fields define the operators:

$SO(1,3)$, the symmetry in $d=(1+3)$ originating in $SO(1,7)$

$\vec{N}_{\pm} = \frac{1}{2}(S^{23} \pm i S^{01}, S^{31} \pm i S^{02}, S^{12} \pm i S^{03})$,

$\vec{\tilde{N}}_{\pm} = \frac{1}{2}\,(\tilde{S}^{23} \pm i \tilde{S}^{01},\, \tilde{S}^{31} \pm i \tilde{S}^{02},\, \tilde{S}^{12} \pm i \tilde{S}^{03})$,

$SU(2)\times SU(2)$, that is of $SO(4)$ originating in $SO(1,7)$

$\vec{\tau}^{1,2} = \frac{1}{2}(S^{58} \mp S^{67},\, S^{57} \pm S^{68},\, S^{56} \mp S^{78})$

$\vec{\tilde{\tau}}^{(1,2)} = \frac{1}{2}(\tilde{S}^{58} \mp \tilde{S}^{67},\, \tilde{S}^{57} \pm \tilde{S}^{68},\, \tilde{S}^{56} \mp \tilde{S}^{78})$

$SU(3)\times U(1)$ originating in $SO(6)$

$\vec{\tau}^{3} := \frac{1}{2}\,\{S^{9\,12} - S^{10\,11},\, S^{9\,11} + S^{10\,12},$
$S^{9\,10} - S^{11\,12}, S^{9\,14} - S^{10\,13},\, S^{9\,13} + S^{10\,14},\, S^{11\,14} - S^{12\,13},$
$S^{11\,13} + S^{12\,14},\, \frac{1}{\sqrt{3}}(S^{9\,10} + S^{11\,12} - 2S^{13\,14})\}$,

$\tau^{4} := -\frac{1}{3}(S^{9\,10} + S^{11\,12} + S^{13\,14})$,

and equivalently in the $\tilde{S}^{ab}$ sector.

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

$$S^{03}, S^{12}, S^{56}, S^{78}, S^{9\,10}, S^{11\,12}, S^{13\,14},$$
$$\tilde{S}^{03}, \tilde{S}^{12}, \tilde{S}^{56}, \tilde{S}^{78}, \tilde{S}^{9\,10}, \tilde{S}^{11\,12}, \tilde{S}^{13\,14}.$$

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

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

S^{ab} generate **the members of one family**. 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 $\tau^{41}=1/6$)

i		$\vert^a\psi_i>$	$\Gamma^{(1,3)}$	S^{12}	$\Gamma^{(4)}$	τ^{13}	τ^{21}	Y	Y'
		Octet, $\Gamma^{(1,7)}=1$, $\Gamma^{(6)}=-1$, of quarks							
1	$\mathbf{u_R^{c1}}$	$\stackrel{03}{(+i)}\stackrel{12}{(+)} \mid \stackrel{56}{(+)}\stackrel{78}{(+)} \mid\mid \stackrel{9\;10}{(+)}\;\stackrel{11\;12}{(-)}\;\stackrel{13\;14}{(-)}$	1	$\frac{1}{2}$	1	0	$\frac{1}{2}$	$\frac{2}{3}$	$-\frac{1}{3}$
2	u_R^{c1}	$\stackrel{03}{[-i]}\stackrel{12}{[-]} \mid \stackrel{56}{(+)}\stackrel{78}{(+)} \mid\mid \stackrel{9\;10}{(+)}\;\stackrel{11\;12}{(-)}\;\stackrel{13\;14}{(-)}$	1	$-\frac{1}{2}$	1	0	$\frac{1}{2}$	$\frac{2}{3}$	$-\frac{1}{3}$
3	d_R^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{(+)} \mid \stackrel{56}{[-]}\stackrel{78}{[-]} \mid\mid \stackrel{9\;10}{(+)}\;\stackrel{11\;12}{(-)}\;\stackrel{13\;14}{(-)}$	1	$\frac{1}{2}$	1	0	$-\frac{1}{2}$	$-\frac{1}{3}$	$\frac{2}{3}$
4	$\mathbf{d_R^{c1}}$	$\stackrel{03}{[-i]}\stackrel{12}{[-]} \mid \stackrel{56}{[-]}\stackrel{78}{[-]} \mid\mid \stackrel{9\;10}{(+)}\;\stackrel{11\;12}{(-)}\;\stackrel{13\;14}{(-)}$	1	$-\frac{1}{2}$	1	0	$-\frac{1}{2}$	$-\frac{1}{3}$	$\frac{2}{3}$
5	d_L^{c1}	$\stackrel{03}{[-i]}\stackrel{12}{(+)} \mid \stackrel{56}{[-]}\stackrel{78}{(+)} \mid\mid \stackrel{9\;10}{(+)}\;\stackrel{11\;12}{(-)}\;\stackrel{13\;14}{(-)}$	-1	$\frac{1}{2}$	-1	$-\frac{1}{2}$	0	$\frac{1}{6}$	$\frac{1}{6}$
6	$\mathbf{d_L^{c1}}$	$\stackrel{03}{(+i)}\stackrel{12}{[-]} \mid \stackrel{56}{[-]}\stackrel{78}{(+)} \mid\mid \stackrel{9\;10}{(+)}\;\stackrel{11\;12}{(-)}\;\stackrel{13\;14}{(-)}$	-1	$-\frac{1}{2}$	-1	$-\frac{1}{2}$	0	$\frac{1}{6}$	$\frac{1}{6}$
7	$\mathbf{u_L^{c1}}$	$\stackrel{03}{[-i]}\stackrel{12}{(+)} \mid \stackrel{56}{(+)}\stackrel{78}{[-]} \mid\mid \stackrel{9\;10}{(+)}\;\stackrel{11\;12}{(-)}\;\stackrel{13\;14}{(-)}$	-1	$\frac{1}{2}$	-1	$\frac{1}{2}$	0	$\frac{1}{6}$	$\frac{1}{6}$
8	u_L^{c1}	$\stackrel{03}{(+i)}\stackrel{12}{[-]} \mid \stackrel{56}{(+)}\stackrel{78}{[-]} \mid\mid \stackrel{9\;10}{(+)}\;\stackrel{11\;12}{(-)}\;\stackrel{13\;14}{(-)}$	-1	$-\frac{1}{2}$	-1	$\frac{1}{2}$	0	$\frac{1}{6}$	$\frac{1}{6}$

$-\mathcal{L}_Y = \psi^\dagger\, \gamma^0 \{\stackrel{78}{(+)}\, p_{0+} + \stackrel{78}{(-)}\, p_{0-}\}\psi$, $\gamma^0\, \stackrel{78}{(-)}$ transforms $\mathbf{u_R}$ of the 1^{st} row into $\mathbf{u_L}$ of the 7^{th} row, while $\gamma^0\, \stackrel{78}{(+)}$ transforms $\mathbf{d_R}$ of the 4^{rd} row into $\mathbf{d_L}$ of the 6^{th} row, doing what the Higgs and γ^0 do in the Stan. model.

$\tilde{\mathbf{S}}^{ab}$ **generate families**. **Both vectors bellow** ,
the second following from the first after the application of $\tilde{S}^{01}$,
describe a **right handed *u*-quark of the same spin and colour**.

$$
\begin{array}{l}
\tilde{S}^{01} \;\; \text{applied} \;\; \text{on} \\
\overset{03}{(+i)}\overset{12}{(+)} \mid \overset{56}{(+)}\overset{78}{(+)} \mid\mid \overset{910}{(+)}\overset{1112}{(-)}\overset{1314}{(-)} \quad \text{generates} \\
\overset{03}{[+i]}\overset{12}{[+]} \mid \overset{56}{(+)}\overset{78}{(+)} \mid\mid \overset{910}{(+)}\overset{1112}{(-)}\overset{1314}{(-)} \,.
\end{array}
$$

Eight families of u_R with the spin $1/2$ of a particular colour and of a **colourless** ν_R :

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	$\stackrel{03}{[+i]}\stackrel{12}{[+]} \vert \stackrel{56}{[+]}\stackrel{78}{[+]} \Vert \stackrel{9\,10}{(+)}\;\stackrel{11\,12}{(+)}\;\stackrel{13\,14}{(+)}$

Before the break of $SO(1,3)\times$ $\mathbf{SU(2)}$ $\times$ $\mathbf{SU(2)}$ $\times U(1) \times SU(3)$ into $SO(1,3)\times$ $\mathbf{SU(2)}$ $\times$ $U(1) \times SU(3)$ all the eight families are massless.

The quantum numbers of the eight families, the same for each member of a particular family

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

The quantum numbers of the members – quarks and leptons, left and right handed – of any of the eight families ($i \in \{1, \cdots, 8\}$) presented above

	$\Gamma^{(1+3)}$	S_L^{03}	S_L^{12}	S_R^{03}	S_R^{12}	τ^{13}	τ^{23}	Y	Q	$SU(3)$	Y'
u_{Li}	-1	$\mp\frac{i}{2}$	$\pm\frac{1}{2}$	0	0	$\frac{1}{2}$	0	$\frac{1}{6}$	$\frac{2}{3}$	triplet	$-\frac{1}{2}\tan^2\theta_2$
d_{Li}	-1	$\mp\frac{i}{2}$	$\pm\frac{1}{2}$	0	0	$-\frac{1}{2}$	0	$\frac{1}{6}$	$-\frac{1}{3}$	triplet	$-\frac{1}{2}\tan^2\theta_2$
ν_{Li}	-1	$\mp\frac{i}{2}$	$\pm\frac{1}{2}$	0	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	0	singlet	$\frac{1}{2}\tan^2\theta_2$
e_{Li}	-1	$\mp\frac{i}{2}$	$\pm\frac{1}{2}$	0	0	$-\frac{1}{2}$	0	$-\frac{1}{2}$	-1	singlet	$\frac{1}{2}\tan^2\theta_2$
u_{Ri}	1	0	0	$\pm\frac{i}{2}$	$\pm\frac{1}{2}$	0	$\frac{1}{2}$	$\frac{2}{3}$	$\frac{2}{3}$	triplet	$\frac{1}{2}(1-\frac{1}{3}\tan^2\theta_2)$
d_{Ri}	1	0	0	$\pm\frac{i}{2}$	$\pm\frac{1}{2}$	0	$-\frac{1}{2}$	$-\frac{1}{3}$	$-\frac{1}{3}$	triplet	$-\frac{1}{2}(1+\frac{1}{3}\tan^2\theta_2)$
ν_{Ri}	1	0	0	$\pm\frac{i}{2}$	$\pm\frac{1}{2}$	0	$\frac{1}{2}$	0	0	singlet	$\frac{1}{2}(1+\frac{1}{3}\tan^2\theta_2)$
e_{Ri}	1	0	0	$\pm\frac{i}{2}$	$\pm\frac{1}{2}$	0	$-\frac{1}{2}$	-1	-1	singlet	$-\frac{1}{2}(1-\frac{1}{3}\tan^2\theta_2)$

BREAKING THE STARTING SYMMETRY SO(1,13)

Breaks of symmetries when starting with **massless spinors** (fermions) and **vielbeins and two kinds of spin connections**

$\mathbf{SO(1,13)}$

BREAK I

$\downarrow$

$\mathbf{SO(1,7)} \times \quad \mathbf{U(1)} \times \quad \mathbf{SU(3)}$

$\swarrow \qquad \searrow$

eight massless families

$\mathbf{SO(1,3)} \times \mathbf{SO(4)} \times \mathbf{U(1)} \times$ (eight massless families) $\mathbf{SU(3)}$

BREAK IIa

$$\underbrace{\mathbf{SO(1,3)} \times \mathbf{SU(2)} \times \mathbf{U(1)} \times \quad (\text{four families stay massless,}) \quad (\text{four families obtain masses}) \quad \mathbf{SU(3)}}$$

$\downarrow$

BREAK IIb

The Standard Model type of breaking

$\downarrow$

$\mathbf{SO(1,3)} \times \mathbf{U(1)} \times \mathbf{SU(3)}$

BREAK I

- Breaking the symmetry from $SO(1,13)$ to $SO(1,7) \times U(1) \times SU(3)$ occurs at very high energy scale (E $> 10^{16}$ GeV). It is followed by the break of $SO(1,7) \times U(1) \times SU(3)$ into $SO(1,3) \times SU(2) \times SU(2) \times U(1) \times SU(3)$. Both breaks leave eight families ($2^{8/2-1} = 8$, determined by the symmetry of $SO(1,7)$) massless.

- We are studying (with H.B. Nielsen, D. Lukman) on a toy model of $d = 1+5$ how to obtain after breaking symmetries massless spinors chirally coupled to the Kaluza-Klein-like gauge fields. Boundaries and the "effective two dimensionalities" seems to be very promising.

The BREAK II a

from $SO(1,3) \times SU(2) \times SU(2) \times U(1) \times SU(3)$ to $SO(1,3) \times SU(2) \times U(1) \times SU(3)$ (at E around 10^{13} GeV or lower).

- The symmetries: **hyper charge ($U(1)$), weak charge ($SU(2)$), spin ($SO(1,3)$)** and **colour ($SU(3)$)** with the corresponding **gauge fields**.
- It leaves **four of the families massless and mass protected** (since they are singlets with respect to one of the two $\tilde{\mathbf{SU}}(\mathbf{2})$ and with only the left handed members carrying the weak charge), the **upper four families obtain masses**.
- There are the vielbeins (${e^a}_\alpha$), together with the two kinds of the spin connection fields, which cause the breaking and bring masses to the spinors and the gauge fields.

The BREAK II b

The break from $SO(1,3) \times SU(2) \times U(1) \times SU(3)$ to $SO(1,3) \times U(1) \times SU(3)$ (the weak scale break) **leads to the massive four lowest families**, three of them are the measured ones, **the fourth is predicted to be possibly seen at the LHC** and to the masses of the **weak bosons**.

- The symmetries: **elm charge ($U(1)$), spin ($SO(1,3)$)** and **colour ($SU(3)$)** charge with the corresponding **gauge fields**.
- **All the members of the eight families have the same family content: the same quarks and leptons, coupling to the same gauge fields,**
  they differ in the masses and mixing matrices.
- Again, the vielbeins, and the two kinds of the spin connection fields cause this weak scale breaking.

The gauge fields, the mass matrices and the scalars after the BREAK II a in the "approach".

The gauge **vector fields** and the **scalars** after the break from $SO(1,3)\times SU(2)\times SU(2)\times U(1)\times SU(3)$ to $SO(1,3)\times SU(2)\times U(1)\times SU(3)$

$$
\begin{aligned}
A^{23}_m &= A^{Y}_m \sin\theta_2 + A^{Y'}_m \cos\theta_2, \quad A^{4}_m = A^{Y}_m \cos\theta_2 - A^{Y'}_m \sin\theta_2,\\
A^{2\pm}_m &= \frac{1}{\sqrt{2}}(A^{21}_m \mp i A^{22}_m),\\
\tilde{A}^{23}_s &= \tilde{A}^{\tilde{Y}}_s \sin\tilde{\theta}_2 + \tilde{A}^{\tilde{Y}'}_s \cos\tilde{\theta}_2, \quad \tilde{A}^{4}_s = \tilde{A}^{\tilde{Y}}_s \cos\tilde{\theta}_2 - \tilde{A}^{\tilde{Y}'}_s \sin\tilde{\theta}_2,\\
\tilde{A}^{2\pm}_s &= \frac{1}{\sqrt{2}}(\tilde{A}^{21}_s \mp i \tilde{A}^{22}_s),
\end{aligned}
$$

for $m=0,1,2,3,\ s=7,8$ and particular values of θ_2 and $\tilde{\theta}_2$. The scalar fields ($\tilde{A}^{\tilde{Y}'}_s$, $\tilde{A}^{2\pm}_s$, $\tilde{A}^{\tilde{N}_R i}_s$, $i=1,2,3,$) are assumed to have nonzero vacuum expectation values.

The **mass matrices** for **quarks and leptons in the $\tilde{S}^{ab}$** sector after the **break of** $SO(1,3) \times SU(2) \times SU(2) \times U(1) \times SU(3)$ **into** $SO(1,3) \times SU(2) \times U(1) \times SU(3)$ on a **tree level**.

	I	*II*	*III*	*IV*	*V*	*VI*	*VII*	*VIII*
I	**0**	**0**	**0**	**0**	0	0	0	0
II	**0**	**0**	**0**	**0**	0	0	0	0
III	**0**	**0**	**0**	**0**	0	0	0	0
IV	**0**	**0**	**0**	**0**	0	0	0	0
V	0	0	0	0	$\frac{1}{2}(\tilde{a}^{23}_{\pm} + \tilde{a}^{\tilde{N}^3_R}_{\pm})$	$-\tilde{\mathbf{a}}^{\mathbf{2}-}_{\pm}$	$-\tilde{\mathbf{a}}^{\tilde{\mathbf{N}}^{+}_{\mathbf{R}}}_{\pm}$	**0**
VI	0	0	0	0	$-\tilde{\mathbf{a}}^{\mathbf{2}+}_{\pm}$	$\frac{1}{2}(-\tilde{a}^{23}_{\pm} + \tilde{a}^{\tilde{N}^3_R}_{\pm})$	**0**	$-\tilde{\mathbf{a}}^{\tilde{\mathbf{N}}^{+}_{\mathbf{R}}}_{\pm}$
VII	0	0	0	0	$-\tilde{\mathbf{a}}^{\tilde{\mathbf{N}}^{-}_{\mathbf{R}}}_{\pm}$	**0**	$\frac{1}{2}(\tilde{a}^{23}_{\pm} - \tilde{a}^{\tilde{N}^3_R}_{\pm})$	$-\tilde{\mathbf{a}}^{\mathbf{2}-}_{\pm}$
VIII	0	0	0	0	**0**	$-\tilde{\mathbf{a}}^{\tilde{\mathbf{N}}^{-}_{\mathbf{R}}}_{\pm}$	$-\tilde{\mathbf{a}}^{\mathbf{2}+}_{\pm}$	$-\frac{1}{2}(\tilde{a}^{23}_{\pm} + \tilde{a}^{\tilde{N}^3_R}_{\pm})$

It then follows that
the gauge fields ($\mathbf{A_m^{Y'}}$, $\mathbf{A_m^{2\pm}}$) of the charges
$\mathbf{Y'} = \tau^{\mathbf{23}} + \tau^{\mathbf{4}} \tan^2\theta_2$,
$\tau^{\mathbf{2\pm}} = \frac{1}{2}((\mathbf{S^{58}} + \mathbf{S^{67}}) \pm i(\mathbf{S^{57}} - \mathbf{S^{68}}))$,
manifest as massive fields,
while the gauge fields ($\mathbf{A_m^{Y}}$, $\mathbf{A_m^{1i}}$, $i = 1, 2, 3$), of the charges
$\mathbf{Y} = \tau^{\mathbf{4}} - \tau^{\mathbf{23}}$,
$\tau^{\mathbf{1i}} = \frac{1}{2}((\mathbf{S^{58}} - \mathbf{S^{67}}), (\mathbf{S^{57}} + \mathbf{S^{68}}), (\mathbf{S^{56}} - \mathbf{S^{78}}))$
stay massless.
There are nonzero vacuum expectation values of the scalar fields ($\mathbf{\tilde{A}_s^{\tilde{Y}'}}$, $\mathbf{\tilde{A}_s^{2\pm}}$), we take $\tan\tilde{\theta}_2 = 0$, and of ($\mathbf{\tilde{A}_s^{\tilde{N}_R i}}$, $i = 1, 2, 3,$), with the scalar vielbeins included, all with respect to $d = (1 + 3)$, which take care on the tree level of the masses of the gauge fields and of the upper four families of quarks and leptons.

- The properties of the upper massive four families are now under consideration. We are treating the one loop and possibly also the two loop corrections to the three level, trying to understand, how do the loop corrections influence the mass matrices when starting from the tree level as presented above. Bellow the tree level the contribution from the massive gauge fields ($\mathbf{A}_{\mathbf{m}}^{\mathbf{Y}'}, \mathbf{A}_{\mathbf{m}}^{\mathbf{2}\pm}$) and the dynamical massive scalar fields ($\tilde{\tilde{\mathbf{A}}}_{\mathbf{s}}^{\mathbf{2}\pm}, \tilde{\mathbf{A}}_{\mathbf{s}}^{\tilde{\mathbf{Y}}'}$), as well as from the dynamical massive scalar fields interacting with the fermions through the operators ($\tilde{N}_R^{\mp}$, $\tilde{N}_R^3$, $\tilde{Y}'$), are studied.
- **The lowest one of the upper four families (all with no Yukawa couplings to the lowest four families) is predicted to constitute the dark matter**.

The scalar, gauge fields and Yukawa couplings after the BREAK II b in the "approach".

The gauge **vector fields** and the **scalars** after the break from $SO(1,3) \times SU(2) \times U(1) \times SU(3)$ to $SO(1,3) \times U(1) \times SU(3)$:

$$
\begin{aligned}
A^{13}_m &= A_m \sin\theta_1 + Z_m \cos\theta_1,\\
A^{Y}_m &= A_m \cos\theta_1 - Z_m \sin\theta_1,\\
W^{\pm}_m &= \frac{1}{\sqrt{2}}(A^{11}_m \mp i A^{12}_m),\\
\tilde{A}^{13}_s &= \tilde{A}^{\tilde{Q}}_s \sin\tilde{\theta}_1 + \tilde{A}^{\tilde{Q}'}_s \cos\tilde{\theta}_1,\\
\tilde{A}^{Y}_s &= \tilde{A}^{\tilde{Q}}_s \cos\tilde{\theta}_1 - \tilde{A}^{\tilde{Q}'}_s \sin\tilde{\theta}_1,\\
\tilde{A}^{1\pm}_s &= \frac{1}{\sqrt{2}}(\tilde{A}^{11}_s \mp i\tilde{A}^{12}_s),
\end{aligned}
$$

for $m = 0,1,2,3$, $s = 7,8$, and particular values of $\theta_1 = \theta_w$ and $\tilde{\theta}_1$. The scalar fields ($A^{Q'}_s$, $A^{1\pm}_s$, and $\tilde{A}^{\tilde{Q}'}_s$, $\tilde{A}^{1\pm}_s$, $\tilde{A}^{\tilde{N}_L i}_s$, $i = 1,2,3,$) are assumed to have nonzero vacuum expectation values.

The **mass matrices** for the lower four families of the **quarks and leptons in the $\tilde{S}^{ab}$** sector after the **break of** $SO(1,3) \times SU(2) \times U(1) \times SU(3)$ **into** $SO(1,3) \times U(1) \times SU(3)$ on a **tree level**.

	I	II	III	IV
I	$\frac{1}{2}\,(\tilde{a}^{13}_{\pm} + \tilde{a}^{\tilde{N}^3_L}_{\pm})$	$\tilde{a}^{1+}_{\pm}$	$\tilde{a}^{\tilde{N}^+_L}_{\pm}$	0
II	$\tilde{a}^{1-}_{\pm}$	$\frac{1}{2}(-\tilde{a}^{13}_{\pm} + \tilde{a}^{\tilde{N}^3_L}_{\pm})$	0	$\tilde{a}^{\tilde{N}^+_L}_{\pm}$
III	$\tilde{a}^{\tilde{N}^-_L}_{\pm}$	0	$\frac{1}{2}\,(\tilde{a}^{13}_{\pm} - \tilde{a}^{\tilde{N}^3_L}_{\pm})$	$\tilde{a}^{1+}_{\pm}$
IV 0	$\tilde{a}^{\tilde{N}^-_L}_{\pm}$	$\tilde{a}^{1-}_{\pm}$	$-\frac{1}{2}\,(\tilde{a}^{13}_{\pm} + \tilde{a}^{\tilde{N}^3_L}_{\pm})$	

It then follows that the gauge fields
$(\mathbf{A}^{\mathbf{Q}'}_{\mathbf{m}} = \mathbf{Z}_{\mathbf{m}}, \mathbf{A}^{\mathbf{1}\pm}_{\mathbf{m}} = \mathbf{W}^{\pm}_{\mathbf{m}})$
of the charges
$\mathbf{Q}' = \tau^{\mathbf{13}} + \mathbf{Y} \tan^2 \theta_1,$
$\tau^{\mathbf{1}\pm} = \frac{1}{2}((\mathbf{S}^{\mathbf{58}} - \mathbf{S}^{\mathbf{67}}) \pm i(\mathbf{S}^{\mathbf{57}} + \mathbf{S}^{\mathbf{68}})),$
manifest as massive fields,
while the gauge field
$\mathbf{A}^{\mathbf{Q}}_{\mathbf{m}} = \mathbf{A}_{\mathbf{m}}$
of the charge
$\mathbf{Q} = \tau^{\mathbf{4}} + \mathbf{S}^{\mathbf{56}},$
stays massless, in agreement with the Standard model.

There are nonzero vacuum expectation values of the scalar fields
$(\mathbf{A}_s^{\mathbf{Q}'}, \mathbf{A}_s^{\mathbf{Y}'})$,
$(\tilde{\mathbf{A}}_s^{\tilde{\mathbf{Q}}'}, \tilde{\mathbf{A}}_s^{1\pm})$, $(\tilde{\mathbf{A}}_s^{\tilde{\mathbf{N}}_L i}, i = 1, 2, 3,)$
with the scalar vielbeins included,
all with respect to $d = (1 + 3)$,
which take care of the masses of the gauge fields and of the masses of the lower four families of quarks and leptons after this break.
$(\mathbf{A}_s^{\mathbf{Q}'}, \mathbf{A}_s^{\mathbf{Y}'})$ bring to each family member a different diagonal contribution to the mass matrices, the same for all the families, while all the fields $\tilde{\mathbf{A}}_s$ distinguish among the family members.

- The properties of the lower four families were on the tree level studied in the paper New Jour. of Phys. **10** (2008) 093002, under the assumption that going bellow the tree level would take care of the differences between quarks and leptons, in addition to the diagonal contributions on the tree level, which are the same for all the families.

- We are studying now the one loop and possibly also the two loop corrections to the three level contributions, hoping to prove that the loop corrections really change the mass matrices in the right direction, that is in agreement with the experimental data. Bellow the tree level the contributions from the massive gauge fields ($\mathbf{Z_m}, \mathbf{W_m^{\pm}}$) and the dynamical massive scalar fields ($\tilde{\tilde{\mathbf{A}}}_{\mathbf{s}}^{\mathbf{1}\pm}, \tilde{\mathbf{A}}_{\mathbf{s}}^{\tilde{\mathbf{Q}}'}$) are studied, as well as from the dynamical massive scalar fields interacting with the fermions through the operators ($\tilde{N}_R^{\pm}$, $\tilde{N}_R^3$).
- The rough estimations done so far predict that the **fourth family might be seen at the LHC or at somewhat higher energies.**

ESTIMATIONS OF THE PROPERTIES FOR THE LOWER FOUR FAMILIES

The **second break influences the massless** (lower four) **families** leading to the masses and mixing matrices of the **three known families** and predicting masses and matrix elements of the **not yet observed fourth family**.

We present here the two years old results for the lowest four families, evaluating the mass matrices on the tree level, under the assumption that while the symmetries of matrices when going bellow the tree level will not change considerably, which means that the matrix elements will approximately keep their relations, the mass matrices of the family members would manifest the differences in all the matrix elements. We accordingly assumed that the mass matrices of the u-quarks, **d**-quarks, ν and **e** are all different.

The mass matrix for the lower four families of **u**-quarks ($-$) and **d**-quarks ($+$) is not assumed to be real and symmetric. Adding to the contributions of the scalar fields originating in $\tilde{sts'}$ also those of ω_{abs}, which distingush among the members of a family, and exchanging the third and the fourth basic state, we parameterize

$$\begin{pmatrix} a_\pm & b_\pm & -c_\pm & 0 \\ b_\pm & a_\pm + d_{1\pm} & 0 & -c_\pm \\ c_\pm & 0 & a_\pm + d_{2\pm} & b_\pm \\ 0 & c_\pm & b_\pm & a_\pm + d_{3\pm} \end{pmatrix}$$

Fitting these parameters with the Monte-Carlo program to the experimental data within the known accuracy and to the **assumed values for the fourth family masses** we get as follows in the next page.

For the **u**-quarks the mass matrix is as follows

$$\begin{pmatrix} (9,22) & (-150,-83) & 0 & (-306,304) \\ (-150,-83) & (1211,1245) & (-306,304) & 0 \\ 0 & (-306,304) & (171600,176400) & (-150,-83) \\ (-306,304) & 0 & (-150,-83) & 200000 \\ \end{pmatrix}$$

and for the **d**-quarks the mass matrix is

$$\begin{pmatrix} (5,11) & (8.2,14.5) & 0 & (174,198) \\ (8.2,14.5) & (83,115) & (174,198) & 0 \\ 0 & (174,198) & (4260,4660) & (8.2,14.5) \\ (174,198) & 0 & (8.2,14.5) & 200000 \\ \end{pmatrix}.$$

This corresponds to the following values for the masses of the **u** and the **d** quarks (**the fourth family masses are assumed**)

$$\begin{aligned} m_{u_i}/GeV &= (0.005, 1.220, 171., 215.), \\ m_{d_i}/GeV &= (0.008, 0.100, 4.500, 285.), \end{aligned}$$

and the mixing matrix for the quarks

$$\begin{pmatrix} -0.974 & -0.226 & -0.00412 & \mathbf{0.00218} \\ 0.226 & -0.973 & -0.0421 & \mathbf{-0.000207} \\ 0.0055 & -0.0419 & 0.999 & \mathbf{0.00294} \\ \mathbf{0.00215} & \mathbf{0.000414} & \mathbf{-0.00293} & \mathbf{0.999} \end{pmatrix}.$$

THE APPROACH MIGHT HAVE THE ANSWER TO THE QUESTION
WHAT DOES CONSTITUTE THE DARK MATTER

The candidate for the Dark matter constituent must have the following properties:

1. It must be stable in comparison with the age of the Universe.
2. Its density distribution within a galaxy is approximately spherically symmetric and decreases approximately with the second power of the radius of the galaxy.
3. The scattering amplitude of a cluster of constituents with the ordinary matter and among the Dark matter clusters must be very weak in order that the predictions would be in agreement with the observations.
4. The Dark matter constituents and accordingly also the clusters had to have a chance to be formed during the evolution of our Universe so that they agree with the today observed properties of the Universe.

- We study the possibility that the **Dark matter constituents are clusters of the stable fifth family of quarks and leptons**, predicted by the **Approach unifying spins and charges** and due to the Approach having negligible Yukawa couplings to the lower four families (in comparison with the age of the universe).
- There are several candidates for the massive Dark matter constituents in the literature, known as **WIMP**s—weakly interacting massive particles.

- Our Dark matter **fifth family quarks, clustered into baryons**, are **are not WIMPS**. They interact during forming clusters in the evolution of the universe with the colour force, the fifth family baryons interact with the **"nuclear like fifth family force"** and in the case that they are very massive the weak force dominates, the fifth family neutrinos interact with the weak force.

What do we know about the properties of the fifth family members so far? (*Phys. Rev. D* **80**, 083534 (2009))

- The masses of the fifth family members lie much above the known three and the predicted fourth family masses—at around 10 **TeV** or higher—and **much bellow** the break of $SO(1,7)$ to $SO(1,3) \times SU(2) \times SU(2)$, which occurs bellow 10^{13} TeV.
- They interact with the **weak, colour and** $U(1)$ interaction.
- When following the fifth family members through the evolution of the universe up to the today's Dark matter several breaks of symmetries and phase transitions occur.
- Knowing their interactions with the gauge fields, we should be able to estimate their interaction with the ordinary matter in direct measurements.

With respect to the discussions above about the properties of the upper four families, it seems meaningful to expect that the members of one family have approximately the same masses due to the expectation (assumption) that no scalar fields originating in $\omega_{sts'}$ have nonzero vacuum expectation value (the scalar fields $\tilde{\omega}_{abs}$, namely, do distinguish among the family members) although studies of the loop corrections will tell whether this expectation is correct. If this expectation is correct, then $n_5 = u_5 d_5 d_5$ is the lightest fifth family baryon, baryon n_5 $(u_5 d_5 d_5)$ is lighter than the baryon p_5 $(u_5 u_5 d_5)$ due to stronger elm repulsion between two u_5 than between two d_5.

Evaluation of the properties of the fifth family baryons

We use the **Bohr (hydrogen)-like model to estimate the binding energy and the size of the fifth family neutron** $(u_5 d_5 d_5)$, assuming that the differences in masses of the fifth family quarks makes the n_5 stable

$$E_{c_5} \approx -3\,\frac{1}{2}\,\left(\frac{2}{3}\,\alpha_c\right)^2\,\frac{m_{q_5}}{2}c^2, \quad r_{c_5} \approx \frac{\hbar c}{\frac{2}{3}\,\alpha_c\,\frac{m_{q_5}}{2}c^2}. \tag{1}$$

The mass of the cluster is approximately

$$m_{c_5}\, c^2 \approx 3 m_{q_5}\, c^2 (1 - (\frac{1}{3}\,\alpha_c)^2) \tag{2}$$

. (We use the factor of $\frac{2}{3}$ for a two quark pair potential and of $\frac{4}{3}$ for an quark and anti-quark pair potential.)

The **Bohr (hydrogen)-like model gives for the fifth family baryon n_5**

$\frac{m_{q_5} c^2}{\mathrm{TeV}}$	α_c	$\frac{E_{c_5}}{m_{q_5} c^2}$	$\frac{r_{c_5}}{10^{-6}\mathrm{fm}}$	$\frac{\Delta m_{ud} c^2}{\mathrm{GeV}}$
1	0.16	-0.016	$3.2 \cdot 10^3$	0.05
10	0.12	-0.009	$4.2 \cdot 10^2$	0.5
10^2	**0.10**	**-0.006**	**52**	**5**
10^3	0.08	-0.004	6.0	50
10^4	0.07	-0.003	0.7	$5 \cdot 10^2$
10^5	0.06	-0.003	0.08	$5 \cdot 10^3$

Table: m_{q_5} in TeV/c^2 is the assumed fifth family quark mass, α_c is the coupling constant of the colour interaction at $E \approx (-E_{c_5}/3)$ which is the kinetic energy of quarks in the baryon, r_{c5} is the corresponding average radius. Then $\sigma_{c_5} = \pi r_{c_5}^2$ is the corresponding scattering cross section.

- **The nucleon-nucleon cross section is for the fifth family nucleons obviously for many orders of magnitude smaller than for the first family nucleons.**
- The binding energy is of the two orders of magnitude smaller than the mass of a cluster at $m_{q_5} \approx 10$ TeV to 10^6 TeV.

Evolution of the abundance of the fifth family members in the universe:

To estimate the behaviour of our stable heavy family of quarks and anti-quarks in the expanding universe we need to know:

- the masses of our fifth family members,
- their particle—anti-particle asymmetry.

We shall take the **fifth family mass** as the **parameter to be determined from the today's Dark matter density** and assume **no fifth family particle—anti-particle asymmetry**.

Both, the masses and the asymmetry follow from our starting Lagrangean, if (when) we would be able to calculate them. But for heavy enough fifth family baryons this asymmetry is not important.

- In the energy interval we treat, the **one gluon exchange is the dominant contribution up to ≈ 1 GeV** when the colour phase transition starts.
- The fifth family quarks and anti-quarks start to freeze out when the temperature of the plasma falls close to $m_{q_5}\, c^2/k_b$.
- They are forming clusters (bound states) when the temperature falls close to the binding energy of the fifth family baryons.

- When the three quarks or three anti-quarks of the fifth family form a **colourless baryon** (or anti-baryon), they **decouple from the rest of the plasma** due to small scattering cross section manifested by the average radius presented in the above Table, manifesting the **"nuclear force" among the fifth family baryons**.
- The fifth family quarks (or coloured clusters), which survive up to the colour phase transition, deplete at the phase transition (before the first family quarks start to form with them the colourless hadrons), due to their very high mass and binding energy. We made a rough estimation of how and why this is happening. The proof has not yet been done.
- More accurate evaluations are in progress.

To follow the behaviour of the fifth family members in the expanding universe, we need to solve the corresponding coupled Boltzmann equations, for which we need to evaluate:

- Their thermally averaged scattering cross sections (as the function of the temperature) for scattering
  i.a.) into all the relativistic quarks and anti-quarks of lower families ($< \sigma v >_{q\bar{q}}$),
  i.b.) into gluons ($< \sigma v >_{gg}$),
  i.c.) into (annihilating) bound states of a fifth family quark and an anti-quark ($< \sigma v >_{(q\bar{q})_b}$),
  i.d.) into bound states of two fifth family quarks and into the fifth family baryons ($< \sigma v >_{c_5}$)
  (and equivalently into two anti-quarks and into anti-baryons).
- The probability for quarks and anti-quarks of the fifth family to annihilate at the colour phase transition ($Tk_b \approx 1$ GeV).

We solve the Boltzmann equation, which treats **in the expanding universe** the number density of all the fifth family quarks as well as of their baryons as a function of thr temperature T (T $= T(t)$, t is the time parameter).

The fifth family quarks scatter with anti-quark into all the other relativistic quarks and anti-quarks ($< \sigma v >_{q\bar{q}}$) and into gluons ($< \sigma v >_{gg}$).

At the beginning, when the quarks are becoming non-relativistic and start to freeze out, the formation of bound states is negligible.

- The Boltzmann equation for the fifth family quarks n_{q_5} (and equivalently for anti-quarks $n_{\bar{q}_5}$

$$a^{-3}\frac{d(a^3\mathbf{n_{q_5}})}{dt} = <\sigma v>_{q\bar{q}}\ n^{(0)}_{q_5} n^{(0)}_{\bar{q}_5}\left(-\frac{n_{q_5} n_{\bar{q}_5}}{n^{(0)}_{q_5} n^{(0)}_{\bar{q}_5}} + \frac{n_q n_{\bar{q}}}{n^{(0)}_{q} n^{(0)}_{\bar{q}}}\right) + <\sigma v>_{gg}\ n^{(0)}_{q_5} n^{(0)}_{\bar{q}_5}\left(-\frac{n_{q_5} n_{\bar{q}_5}}{n^{(0)}_{q_5} n^{(0)}_{\bar{q}_5}} + \frac{n_g n_g}{n^{(0)}_{g} n^{(0)}_{g}}\right).$$

$n_i^{(0)} = g_i\,(\frac{m_i c^2\, Tk_b}{(\hbar c)^2})^{\frac{3}{2}}\, e^{-\frac{m_i c^2}{Tk_b}}$ for $m_i c^2 >> Tk_b$ (which is our case and to $\frac{g_i}{\pi^2}\,(\frac{Tk_b}{\hbar c})^3$ for $m_i c^2 << Tk_b$).

- When the temperature of the expanding universe falls close enough to the binding energy of the cluster of the fifth family quarks (and anti-quarks), the bound states of quarks (and anti-quarks) and the clusters of fifth family baryons (in our case neutrons n_5) (and anti-baryons $\bar{n}_5$—anti-neutrons) start to be formed.
- The corresponding Boltzmann equation for the number of baryons n_{c_5} reads

$$a^{-3}\,\frac{d(a^3 \mathbf{n_{c_5}})}{dt} \;=\; <\sigma v>_{c_5}\; n^{(0)^2}_{q_5} \left(\left(\frac{\mathbf{n_{q_5}}}{n^{(0)}_{q_5}} \right)^2 - \frac{n_{c_5}}{n^{(0)}_{c_5}} \right).$$

The number density of the fifth family quarks $\mathbf{n_{q_5}}$ ($\bar{\mathbf{n}}_{\mathbf{q_5}}$) which has above the temperature of the binding energy of the clusters of the fifth family quarks (almost) reached the decoupled value, starts to decrease again due to the formation of the clusters.

$$
\begin{aligned}
& a^{-3}\frac{d(a^3 \mathbf{n_{q_5}})}{dt} = \\
& <\sigma v>_{c_5} \; n^{(0)}_{q_5}\, n^{(0)}_{q_5} \left[-\left(\frac{n_{q_5}}{n^{(0)}_{q_5}}\right)^2 + \frac{\mathbf{n_{c_5}}}{n^{(0)}_{c_5}} - \frac{\eta_{(q\bar{q})_b}}{\eta_{c_5}} \left(\frac{n_{q_5}}{n^{(0)}_{q_5}}\right)^2 \right] + \\
& <\sigma v>_{q\bar{q}} \; n^{(0)}_{q_5}\, n^{(0)}_{\bar{q}_5} \left(-\frac{n_{q_5}\, n_{\bar{q}_5}}{n^{(0)}_{q_5}\, n^{(0)}_{\bar{q}_5}} + \frac{n_q\, n_{\bar{q}}}{n^{(0)}_q\, n^{(0)}_{\bar{q}}}\right) + \\
& <\sigma v>_{gg} \; n^{(0)}_{q_5}\, n^{(0)}_{\bar{q}_5} \left(-\frac{n_{q_5}\, n_{\bar{q}_5}}{n^{(0)}_{q_5}\, n^{(0)}_{\bar{q}_5}} + \frac{n_g\, n_g}{n^{(0)}_g\, n^{(0)}_g}\right).
\end{aligned}
$$

- At the temperature < 1 **GeV/c^2 the colour phase transition starts** and **the fifth family quarks and anti-quarks and the coloured fifth family clusters**, all with a very large mass (several 10^8 MeV to be compared with 300 MeV of the first family dressed quarks) and accordingly with the very large momentum, with a very large binding energy (see Table above) and also with the large scattering cross section (which all the quarks obtain at the coloured phase transition) **deplete before forming the hadrons with the lower family members**. This is waiting to be proved
- The colourless **fifth family baryons**, being bound into very small clusters, **do not feel the colour phase transition**.

Evolution

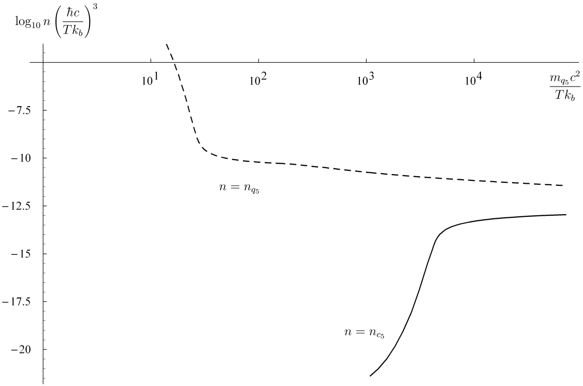

Figure: The dependence of the two number densities n_{q_5} (of the fifth family quarks) and n_{c_5} (of the fifth family clusters) as the function of $\frac{m_{q_5} c^2}{T k_b}$ is presented for the values $m_{q_5} c^2 = 71\ \mathrm{TeV}$, $\eta_{c_5} = \frac{1}{50}$ and $\eta_{(q\bar{q})_b} = 1$. We take $g^* = 91.5$.

From the calculated decoupled number density of baryons and anti-baryons of the fifth family quarks (anti-quarks) $n_{c_5}(T_1)$ at the temperature $T_1 k_b = 1$ GeV, where we stopped our calculations (as a function of the quark mass and of the two parameters η_{c_5} and $\eta_{(q\bar{q})_b}$, which measure the inaccuracy of our calculations), the today's mass density of the dark matter follows $a_1^3 n_{c_5}(T_1) = a_2^3 n_{c_5}(T_2)$, with the today's $a_0 = 1$ and $T_0 = 2.7..^0$ K)

- $$\rho_{\mathbf{dm}} \quad = \quad \Omega_{dm}\rho_{cr} = 2\,\mathbf{m_{c_5}}\,\mathbf{n_{c_5}(T_1)}\,\left(\frac{T_0}{T_1}\right)^3 \frac{g^*(T_1)}{g^*(T_0)},$$

$\frac{m_{q_5} c^2}{\mathrm{TeV}}$	$\eta_{(q\bar{q})_b} = \frac{1}{10}$	$= \frac{1}{3}$	$= 1$	$= 3$	$= 10$
$\eta_{c_5} = \frac{1}{50}$	21	36	71	159	417
$\eta_{c_5} = \frac{1}{10}$	12	20	39	84	**215**
$\eta_{c_5} = \frac{1}{3}$	9	14	25	54	134
$\eta_{c_5} = 1$	8	11	19	37	88
$\eta_{c_5} = 3$	7	**10**	15	27	60
$\eta_{c_5} = 10$	7*	8*	13	22	43

Table: **The fifth family quark mass** is presented, calculated for different choices of η_{c_5} and of $\eta_{(q\bar{q})_b}$, which take care of the inaccuracy of our calculations.

We read from the above Table the **mass interval** for the **fifth family quarks' mass**

$$\mathbf{10}\ \mathrm{TeV} < \mathbf{m_{q_5}\, c^2} < \mathbf{4} \cdot \mathbf{10^2}\mathrm{TeV}. \tag{3}$$

From this mass interval we estimate from the Bohr-like model the **cross section** for the **fifth family neutrons** $\pi(r_{c_5})^2$:

$$\mathbf{10^{-8}}\mathrm{fm}^{\mathbf{2}} < \sigma_{\mathbf{c_5}} < \mathbf{10^{-6}}\mathrm{fm}^{\mathbf{2}}. \tag{4}$$

It is at least $10^{-6}\times$ smaller than the cross section for the first family neutrons.

Dynamics of the heavy family baryons in our galaxy

- Our Sun's velocity: $v_S \approx (170 - 270)$ **km/s.**
- **Locally dark matter density ρ_{dm} is known within a factor of 10** accurately:
  $\rho_{dm} = \rho_0\, \varepsilon_\rho, \rho_0 = 0.3 GeV/(c^2\, cm^3)$,
  we put $\frac{1}{3} < \varepsilon_\rho < 3$.
- The **local velocity of the dark matter clusters $\vec{v}_{dm}$ is unknown**, the estimations are **very model dependant**.
- The velocity of the Earth around the center of the galaxy is equal to: $\vec{v}_E = \vec{v}_S + \vec{v}_{ES}$,
  $v_{ES} = 30$ km/s,
  $\frac{\vec{v}_S \cdot \vec{v}_{ES}}{v_S\, v_{es}} \approx \cos\theta \sin\omega t, \theta = 60^0$.

- **The flux** per unit time and unit surface of our Dark matter clusters hitting the Earth:
  $\Phi_{dm} = \sum_i \frac{\rho_{dmi}}{m_{c_5}} \left|\vec{v}_{dm\,i} - \vec{v}_E\right|$ is $\approx$ equal to

$$\Phi_{dm} \approx \sum_i \frac{\rho_{dmi}}{m_{c_5}} \{\left|\vec{v}_{dm\,i} - \vec{v}_S\right| - \vec{v}_{ES} \cdot \frac{\vec{v}_{dmi} - \vec{v}_S}{\left|\vec{v}_{dmi} - \vec{v}_S\right|}\}.$$

- **We assume** $\sum_i \left|\vec{v_{dmi}} - \vec{v_S}\right| \rho_{dmi} = \varepsilon_{v_{dmS}}\, \varepsilon_\rho\, v_S\, \rho_0$, and correspondingly
- $\sum_i \vec{v}_{ES} \cdot \frac{\vec{v}_{dmi} - \vec{v}_S}{\left|\vec{v}_{dmi} - \vec{v}_S\right|} = v_{ES} \varepsilon_{v_{dmES}} \cos\theta \, \sin\omega t$, with ω for our Earth rotation around our Sun.
- We evaluate $\frac{1}{3} < \varepsilon_{v_{dmS}} < 3$ **and** $\frac{1}{3} < \frac{\varepsilon_{v_{dmS}}}{\varepsilon_{v_{dmES}}} < 3$.

The cross section for our heavy dark matter baryon n_5 to **elastically** scatter on an **ordinary nucleus** with A nucleons in the Born approximation:
$\sigma_{c_5 A} = \frac{1}{\pi \hbar^2} < |M_{c_5 A}| >^2 \; m_A^2$,
$m_A \approx= m_{n_1} \, A^2 \cdots$ the mass of the ordinary nucleus,
$\sigma(A) = \sigma_0 \, A^4$,

- $\sigma_0 = 9 \, \pi r_{c_5}^2 \, \varepsilon_{\sigma_{nucl}}, \quad \frac{1}{30} < \varepsilon_{\sigma_{nucl}} < 30$, when the **"nuclear force"** dominates,
- $\sigma_0 = \frac{m_{n_1} G_F}{\sqrt{2\pi}} (\frac{A-Z}{A})^2 \varepsilon_{\sigma_{weak}} \; (= (10^{-6} fm \frac{A-Z}{A})^2 \, \varepsilon_{\sigma_{weak}})$, $\varepsilon_{\sigma_{weak}} \approx 1$, when the **weak force** dominates ($m_{q_5} > 10^4$ TeV).
- The scattering cross section **among** our heavy neutral baryons n_5 is determined by the weak interaction: $\sigma_{c_5} \approx (10^{-6} \, \mathrm{fm})^2 \, \frac{m_{c_5}}{\mathrm{GeV}}$.

Direct measurements of the fifth family baryons as dark matter constituents:

- Let us assume that DAMA/NaI and CDMS measure our heavy dark matter clusters. We wait for the last analyses in Xe experiment to evaluate their data in the same way.
- **We look for limitations these two experiments might put on properties of our heavy family members.**
- Let an experiment has N_A nuclei per kg with A nucleons.
- At $v_{dmE} \approx 200$ km/s are the $3A$ scatters strongly bound in the nucleus, so that the whole nucleus with A nucleons elastically scatters on a heavy dark matter cluster.
- The number of events per second (R_A) taking place in N_A nuclei is equal to (the cross section is at these energies almost independent of the velocity) what follows

$$\mathbf{R_A = N_A \frac{\rho_0}{m_{c_5}} \sigma(A)\, v_S\, \varepsilon_{v_{dmS}}\, \varepsilon_\rho \left(1 + \frac{\varepsilon_{v_{dmES}}}{\varepsilon_{v_{dmS}}} \frac{v_{ES}}{v_S} \cos\theta \sin \omega t\right)},$$

$\Delta R_A = R_A(\omega t = \frac{\pi}{2}) - R_A(\omega t = 0) = N_A\, R_0\, A^4\, \frac{\varepsilon_{v_{dmES}}}{\varepsilon_{v_{dmS}}} \frac{v_{ES}}{v_S} \cos\theta$,
$R_0 = \sigma_0\, \rho_0 3\, m_{q_5}\, v_S\, \varepsilon$.

$\varepsilon = \varepsilon_\rho\, \varepsilon_{\mathbf{v_{dmES}}}\, \varepsilon_\sigma$,
$\mathbf{10^{-3}} < \varepsilon < \mathbf{10^2}$, for the "nuclear-like force" dominating
$\mathbf{10^{-2}} < \varepsilon < \mathbf{10^1}$, for the weak force dominating

Let $\varepsilon_{\mathbf{cut\,A}}$ determine the efficiency of a particular experiment to detect a dark matter cluster collision, then

$$\mathbf{R_{A\,exp} \approx N_A\, R_0\, A^4\, \varepsilon_{cut\,A} = \Delta R_A \varepsilon_{cut\,A} \frac{\varepsilon_{v_{dmS}}}{\varepsilon_{v_{dmES}}} \frac{v_S}{v_{ES}\ \cos\theta}}.$$

If DAMA/NaI is measuring our heavy family baryons (scattering mostly on I, $A_I = 127$, we neglect Na, with $A = 23$)

$R_{I\,dama} \approx \Delta R_{dama} \frac{\varepsilon_{v_{dmS}}}{\varepsilon_{v_{dmES}}}\,\frac{v_S}{v_{SE}\,\cos 60^0}$,

most of unknowns are hidden in ΔR_{dama}.

For Sun's velocities $v_S = 100, 170, 220, 270$ km/s
we find $\frac{v_S}{v_{SE}\cos\theta} = 7, 10, 14, 18$ respectively.

DAMA/NaI publishes $\Delta R_{I\,dama} = 0,052$ **counts per day and per kg of NaI.**

Then $R_{I\,dama} = 0,052\,\frac{\varepsilon_{v_{dmS}}}{\varepsilon_{v_{dmES}}}\,\frac{v_S}{v_{SE}\cos\theta}$ counts per day and per kg.

CDMS should then in $121 \cdot 2$ days with 1 kg of Ge ($A = 73$) detect

$$R_{Ge}\, \varepsilon_{cut\ cdms} \approx \frac{8.3}{4.0} \left(\frac{73}{127}\right)^4 \frac{\varepsilon_{cut\ cdms}}{\varepsilon_{cut\ dama}} \frac{\varepsilon_{v_{dmS}}}{\varepsilon_{v_{dmES}}} \frac{v_S}{v_{SE} \cos\theta}\, 0.052 \cdot 121 \cdot 2\ ,$$

which is for $v_S = 100, 170, 220, 270$ km/s

equal to $(20,\ 32,\ 42,\ 50)\, \frac{\varepsilon_{cut\ cdms}}{\varepsilon_{cut\ dama}} \frac{\varepsilon_{v_{dmS}}}{\varepsilon_{v_{dmES}}}$.

CDMS has found no event.

If $\frac{\varepsilon_{cut\ cdms}}{\varepsilon_{cut\ dama}} \frac{\varepsilon_{v_{dmS}}}{\varepsilon_{v_{dmES}}}$ is small enough, CDMS will measure our fifth family clusters in the near future.

- **DAMA limits the mass of our fifth family quarks** to $200\,TeV < m_{q_5} c^2 < 10^5\,TeV$.
- **Cosmological evolution requires that masses of the fifth family quarks are not larger than a few 100 TeV.**
- We have checked also all the other cosmological and noncosmological measurements.
- **None** is up to now in contradiction with the prediction of the **Approach unifying spin and charges** that the fifth family members, with neutrino included, constitute the Dark Matter.
- **The Xe experiment has not yet been carefully evaluated.**

- **Also fifth family neutrinos might contribute to the dark matter. We are studying the behaviour of these neutrinos in the expanding universe taking into account annihilation properties of the weakly interacting particles bellow the weak symmetry breaks.**

CONCLUDING REMARKS

There is a lot of open questions which the Standard model of the electroweak and colour interactions leaves unanswered.

The approach unifying spin and charges is offering the (new) way to answer these questions:

- **It offers the explanation of the origin of the charges**, **of the gauge fields and of the scalar fields**.
- **It offers the mechanism for generating families** (the only mechanism in the literature, to my knowledge, which does not on one or another way put the families by hand) and correspondingly **explains the origin of the mass matrices**.

It predicts:

- **In the low energy regime four families (with nonzero Yukawa couplings)**, **the fourth to be possibly seen at the LHC** and correspondingly the masses and the mixing matrices.
- **The stable fifth family which is the candidate to form the dark matter**.

We evaluated:

- **The properties of the lower four families**. We are not yet able to tell the masses of the fourth family members. We are studying their properties bellow the tree level.
- **The properties of the higher four families**, on the tree and bellow the tree level, are under considerations.
- **Possible ways of breaking the starting symmetry, leading to known gauge fields and known properties of quarks and leptons**. We succeeded to guarantee massless spinors chirally coupled through the Kaluza-Klein-like charges to the corresponding gauge fields for a toy model without a fine tuning, which seems to be very promising also for our realistic case.

- **The properties of the fifth family quarks**:
  1. **Their forming the neutral clusters.**
  2. **Their decoupling from the rest of plasma in the evolution of the universe.**
  3. **Their interaction with the ordinary matter (with the first family baryons) and among themselves.**
  4. **Their behaviour in the colour phase transition**. More accurate evaluations are now in process.

I am concluding:

- **There are more than the observed three families**, the **fourth family will possibly be seen at the LHC.**
- **The fifth family, decoupled from the lower four families** (no Yukawa couplings to the lower four families), is the candidate to form the dark matter, **provided that the mass of the fifth family quarks is a few hundred TeV**.
- **I am also predicting, that if DAMA experiments measure our fifth family neutrons, the other direct experiments will "see" the dark matter soon.**

Open problems to be solved—some main steps are already done or are in the process:

- **The way how does the breaking of symmetries occur and define the scales.**
- **The behaviour of quarks and leptons** (neutrinos) **and gauge fields at the phase transitions of the plasma (**$SU(2)$ **and** $SU(3)$**)**.
- **The way how do the loop corrections influence the Yukawa couplings evaluated on the tree level and define correspondingly the differences in masses and mixing matrices.**
- **The discrete symmetries and their nonconservation within the Approach.**
- **Many other not jet solved problems**.