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SIMMETRY BREAKING IN SO(10) GRAND UNIFIED THEORY AND QUARK MASSES

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ՀԱՄԱՉԱՓՈՒԹՅԱՆ ԽԱԽՏՈՒՄԸ ՄԵԾ ՄԻԱՎՈՐՄԱՆ SO/10/ ՄՈԴԵԼՈՒՄ ԵՎ ՔՎԱՐԿՆԵՐԻ ԶԱՆԳՎԱԾՆԵՐԸ

Առաջարկվում է մեծ միավորման SO/10/ մոդել, որտեղ լրացուցիչ ընդհատ համաչափության շնորհիվ լուծված է քվարկների զանգվածների հետ կապված պրոբլեմը, որն առաջանում է SO/10/-ի խախտման այնպիսի սխեմաներում, որոնք օժտված են միջանկյալ $SU(4) \times SU(2)_L \times SU(2)_R$ կամ $SU(3)_c \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$ համաչափությամբ: Հաշվի առնելով այդ ընդհատ համաչափության խախտումը, միայն մեկ թեթև Հիգսի բոզոնի գոյության պայմանից ստացված է b և t քվարկների զանգվածների միջև մի առնչություն, որը թույլ է տալիս որոշել մեծ միավորման M_X մասշտաբի և քվարկ-լեպտոնային համաչափության խախտման M_c մասշտաբի հարաբերությունը: M_X-ի և M_c-ի, ինչպես նաև $SU/2/_R$-համաչափության խախտման M_R մասշտաբի կոնկրետ արժեքները կախված են Վայնբերգի անկյան փորձնական արժեքից և համաձայնվում են պրոտոնի տրոհումը փնտրող փորձերի հետ: Առաջարկված մոդելի սահմաններում ուսումնասիրվում են քվարկների երեք սերունդների խառնման հարցերը: Կոբայաշի-Մասկավայի խառնման անկյունների ստացված արժեքները լավ համաձայնվում են փորձի հետ:

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НАРУШЕНИЕ СИММЕТРИИ В ТЕОРИИ ВЕЛИКОГО ОБЪЕДИНЕНИЯ SU(10) И МАССЫ КВАРКОВ

Предлагается модель великого объединения $SU(10)$, где благодаря дополнительной дискретной симметрии решена проблема с массовым спектром кварков, которая возникает в схемах нарушения $SU(10)$ с промежуточной $SU(4) \times SU(2)_L \times SU(2)_R$ или $SU(3)_c \times U(1)_{B-L} \times SU(2) \times SU(2)_R$ симметрией. С учетом нарушения этой дискретной симметрии и из условия существования только одного легкого бозона Хиггса получено соотношение между массами b и t кварков, которое позволяет зафиксировать отношение масштаба большого объединения M_X и масштаба нарушения кварк-лептонной симметрии M_C. Конкретные значения M_X, M_C, а также масштаба нарушения $SU(2)_R$-симметрии M_R зависят от экспериментального значения угла Вайнберга и согласуются с экспериментами по распаду протона. В рамках предложенной модели изучаются вопросы смешивания трех поколений кварков. Полученные при этом углы смешивания Кобаяши-Маскова хорошо согласуются с опытом.

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SYMMETRY BREAKING IN SO(10) GRAND UNIFIED THEORY AND QUARK MASSES

A model of SO(10) grand unified theory is proposed, where the quark mass spectrum problem occurring in SO(10) breaking schemes with intermediate $SU(4) \times SU(2)_L \times SU(2)_R$ or $SU(3)_c \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$ symmetry is solved owing to additional discrete symmetry. With account of this discrete symmetry breaking and from the condition of existence of only one light Higgs boson, a relation between b and t quark masses is obtained, which enables one to fix up the ratio of grand unification scale M_X to quark-lepton symmetry breaking scale M_C. The concrete values of M_X, M_C as well as $SU(2)_R$-symmetry breaking scale M_R depend on experimental value of Weinberg's angle and agree with the proton decay experiments. The problems of mixing of three generations of quarks are being studied within the proposed model. The obtained Kobayashi-Maskawa mixing angles are in good agreement with experiment.

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Notwithstanding the failure suffered by SU(5) grand unified simplest model in connection with the recent proton decay experiments, the idea of grand unification is still being attractive. The grand unified model based on SO(10) group is the simplest among other possible schemes. This model, just like the SU(5), contains only ordinary quarks and leptons as well as the right neutrino. Whereas in other models, the new particles - mirror fermions interacting with the right currents - are necessarily present.

The question of proton lifetime in the SO(10) was studied in many works [1-4] . It was shown that the proton lifetime larger against one in SU(5) can be obtained only in those SO(10) breaking schemes which contain intermediate symmetries $G = SU(4) \times SU(2)_L \times SU(2)_R$ or $G_1 = SU(3)_c \times U(1)_{B-L} \times$ $\times\, SU(2)_L \times SU(2)_R$. In such schemes it is assumed that $SU(2)_R$ is broken at a considerably lesser scale M_R compared to the grand unification mass M_X , in order to provide an agreement with the proton decay experiments. Whereas available in SO(10) left-right discrete symmetry D, which brings to equality of gauge coupling constants of $SU(2)_L$ and $SU(2)_R$ groups, can be broken at a scale M_P larger compared to M_R [5] , this ensuring an agreement to the cosmology [6,7] . However for such models there arises a problem with fermion masses, which was pointed out by Svetovoy [8]. He has shown that in models with intermediate $SU(2)_L \times SU(2)_R$-symmetry under assumption of existence of only one light (with mass $\sim M_W$) Higgs boson, the masses of b and t quarks turn out to be the same. In our work we propose the SO(10) grand unified model where this difficulty is avoided by means of introducing the additional discrete symmetry. Our obtained relation between b and t quark masses dictates a certain hierarchy in SO(10) breaking scales, when left-right discrete symmetry breaking scale M_P and quark-lepton symmetry breaking scale M_c are not much less than M_X (~ 10 times), while the $SU(2)_R$ breaking scale M_R may be much less than M_X. The proton lifetime predicted may be considerably larger than experimental limit: $\tau_{p \to e^+\pi^0} > 2\cdot 10^{32}$ years. Good predictions for Kobayashi-Maskawa mixing angles are also obtained.

First, observe briefly the arguments given in Ref. [8]. Suppose that at SO(10) breaking there arises an intermediate symmetry G or G_1. For simplicity, we shall assume that all quarks and leptons acquire mass due to interaction with vacuum expectation value (v.e.v.) of one complex Higgs field in the $\underline{10}$ representation (this does not affect the generality of conclusions). The decomposition of this representation of SO(10) under the group G has the form:

$$\underline{10} = (6, 1, 1) + (1, 2, 2) \tag{1}$$

Denote $\Phi = (1, 2, 2)$; precisely this field receives v.e.v. connected with $SU(2)_L$-symmetry breaking. Φ may be presented as a matrix 2 x 2 :

$$\Phi = \begin{pmatrix} \eta^0 & \xi^+ \\ \eta^- & -\xi^{0*} \end{pmatrix} \tag{2}$$

where $\xi = \begin{pmatrix} \xi^0 \\ \xi^- \end{pmatrix}$ and $\eta = \begin{pmatrix} \eta^0 \\ \eta^- \end{pmatrix}$ form $SU(2)_L$ doublets. For simplicity, we can determine also field $\tilde{\Phi}$ conjugated to Φ :

$$\tilde{\Phi} = \begin{pmatrix} \xi^0 & \eta^+ \\ \xi^- & -\eta^{0*} \end{pmatrix} \tag{3}$$

The mass term in the Lagrangian for Φ , $\tilde{\Phi}$ fields has the form:

$$\mathcal{L}_M = \alpha M_1^2 (\Phi^+\Phi + \tilde{\Phi}^+\tilde{\Phi}) + \beta M_2^2 (\Phi^+\tilde{\Phi} + \tilde{\Phi}^+\Phi) \tag{4}$$

where M_1 , M_2 , generally speaking, are of the order of grand unification mass M_X . Diagonal are the states:

$$\Phi_{1,2} = \frac{1}{\sqrt{2}} (\Phi \pm \tilde{\Phi}) \tag{5}$$

with masses $\alpha M_1^2 \pm \beta M_2^2$. One of these states is the Higgs field, whose v.e.v. breaks down the $SU(2)_L \times U(1)$ -symmetry and gives mass to W, Z bosons and fermions. The fine-tuning condition requires that a mass of this field (e.g., Φ_1) should be of M_W order. According to the assumption on the existence of only one light Higgs boson, the other field (Φ_2) must have a mass of M_X order and zero v.e.v. in accordance with the generalized survival hypothesis [9] . This gives equality of η^0 and ξ^0 fields v.e.v. $\langle\eta^0\rangle = \langle\xi^0\rangle$. The fermions obtain mass from Φ and $\tilde{\Phi}$ couplings:

$$a \bar{\Psi}_L \Phi \Psi_R + b \bar{\Psi}_L \tilde{\Phi} \Psi_R + \text{c.c.} \tag{6}$$

With account of the equality of the η^0 and ξ^0 v.e.v., from (6) we obtain for the b and t quarks doublet $\Psi_{L,R} = \begin{pmatrix} t \\ b \end{pmatrix}_{L,R}$ the mass equality $m_b = m_t$, this contradicting the experiment. This relation holds if we add several scalar <u>10</u>-plets or <u>126</u>-plets interacting with the fermion <u>16</u>-plet, only to save the condition of existence of only one light Higgs field. The relation $m_b = m_t$ may be violated only by corrections $\sim M_R/M_X$; however these corrections are small: for the proton lifetime to be large enough, M_R must differ considerably from M_X [2] .

To solve this problem, the author of [10] suggested to introduce additional heavy fermions, this allowing one to avoid the undesirable equality of b and t quark masses.

In this work we suggest a model where this problem is solved in a more economic and natural way. Here we preserve the assumption [8] on the same order of coupling constants and also assume that the generalized survival hypothesis as well as minimal number of fine-tuning conditions take place. Compared to the usual SO(10) model, our model contains only an additional discrete symmetry. Under the action of this discrete symmetry, the Higgs field <u>10</u> and <u>16</u>-plet fermions are transformed as follows:

$$\underline{10} \longrightarrow i\,\underline{10} , \qquad \underline{16} \longrightarrow e^{-i\frac{\pi}{4}}\,\underline{16} \qquad (7)$$

Such symmetry brings to the fact that in mass term (4) $\beta = 0$, and the mass matrix is diagonal in $\phi , \tilde{\phi}$ space. If taking into account also $SU(2)_R$ -symmetry breaking, then it is more suitable to write the mass term in terms of η and ξ spinors:

$$\mathcal{L}_M' = (\alpha M_1^2 + \delta M_R^2)\,\eta^*\eta + (\alpha M_1^2 - \delta M_R^2)\,\xi^*\xi \qquad (8)$$

The fine-tuning condition requires that one of spinors η , ξ , say, η , should have mass $\sim M_W$; then the ξ mass will be $\sim M_R$ and $\langle \xi^\circ \rangle = 0$. Finally we'll obtain that the b quark mass is zero since the second term in (6) is forbidden due to discrete symmetry (7).

One would think, that by introducing discrete symmetry we merely replace one wrong mass relation by another, but that **is not** really the case. One should take into account the discrete symmetry breaking, due to which non-diagonal in η , ξ terms must arise in the mass term, and just then the mass of b quark arises.

Such breaking of discrete symmetry may arise owing to radiative corrections, if taking into account the interaction of the $\underline{10}$ Higgs field with other scalar fields connected with SO(10) breaking: $\underline{54}$, $\underline{210}$, $\underline{45}$, $\underline{126}$. These fields, affected by discrete symmetry, are transformed as follows:

$$\begin{aligned} &\underline{54} \longrightarrow \underline{54} \,, \qquad \underline{210} \longrightarrow \underline{210} \\ &\underline{45} \longrightarrow -\underline{45} \,, \qquad \underline{126} \longrightarrow -\underline{126} \end{aligned} \tag{9}$$

The corresponding breaking scheme looks as

$$\begin{aligned} SO(10) &\xrightarrow{\langle (1,\,1,\,1)_{54}\rangle \,=\, M_X} SU(4) \times SU(2)_L \times SU(2)_R \times D \\ &\xrightarrow{\langle (1,\,1,\,1)_{210}\rangle \,=\, M_P} SU(4) \times SU(2)_L \times SU(2)_R \\ &\xrightarrow{\langle (15,\,1,\,1)_{45}\rangle \,=\, M_C} SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R \\ &\xrightarrow{\langle (\overline{10},\,1,\,3)_{126}\rangle \,=\, M_R} SU(3)_C \times SU(2)_L \times U(1)_y = G_0 \end{aligned} \tag{10}$$

where $(1, 1, 1)_{54}$, $(1, 1, 1)_{210}$, $(15, 1, 1)_{45}$, $(\overline{10}, 1, 3)_{126}$ respectively enter into the decomposition of representations $\underline{54}$, $\underline{210}$, $\underline{45}$, $\underline{126}$ under the group G . The v.e.v. of the $\underline{45}$, $\underline{126}$ fields simultaneously violate also the discrete symmetry (7), (9) introduced by us. Note that with account of (10), the term in (8) proportional to M_R^2 arises from the interaction of

$\underline{10}^*$ $\underline{10}$ $\underline{126}$ $\underline{126}$.

Nondiagonal terms in the mass matrix of η , ξ may arise from the quartic interactions of Higgs field $\underline{10}$ with itself and other fields. Due to our introduced discrete symmetry, such interactions can have only the following form: $\underline{10}\ \underline{10}\ \underline{45}\ \underline{54}$, $\underline{10}\ \underline{10}\ \underline{126}\ \underline{126}$, $\underline{10}\ \underline{10}\ \overline{\underline{126}}\ \overline{\underline{126}}$ plus complex-conjugated terms. However, a concrete group analysis shows that a part of field $\underline{10}$, $\Phi = (1, 2, 2)$, does not obtain mass from these interactions, when fields $\underline{45}$, $\underline{54}$, $\underline{126}$ obtain v.e.v. according to (10). For example, for the first type of interaction this follows from the fact that the product $(1, 2, 2)_{10} \times (1, 2, 2)_{10} \times (15, 1, 1)_{45} \times (1, 1, 1)_{54}$ contains no singlet under the group G.

Stated above implies that a quadratic in Φ term in the mass matrix may arise only after account of radiative corrections. Analysis of all possible types of interactions shows that the effective interactions of the type of $\frac{1}{M_X^2}\ \underline{10}\ \underline{10}\ \underline{45}\ \underline{54}\ \underline{45}\ \underline{45}$ and others not containing field $\underline{126}$ do not give a mass term for Φ . A simplest effective interaction giving rise to such term has the form: $\frac{1}{M_X^2}\ \underline{10}\ \underline{10}\ \underline{45}\ \underline{54}\ \overline{\underline{126}}\ \underline{126}$. Such interaction may arise from diagrams shown in Fig. 1. Fig. 2 presents the same diagrams, but in terms of scalar field representations already decomposed under the G group.

As one can see from Fig. 2, all fields passing along the inner lines of the diagrams must have masses $\sim M_X$. For estimates, we'll consider these masses nearly the same; coupling constants in different vertices of diagrams we'll consider equal too. Finally, for nondiagonal in η , ξ mass term we obtain (contributions of Fig. 1a and 1b diagrams are of the same order):

$$\Delta \mathcal{L}_M = \lambda \gamma \frac{M_c}{M_X} M_R^2 \left(\eta^* \xi + \xi^* \eta \right) \tag{11}$$

where λ is a typical coupling constant.

With account of (8) we obtain the following mass matrix for scalar $SU(2)_L$ -doublets η, ξ :

$$\begin{pmatrix} \alpha M_1^2 + \delta M_R^2 & \lambda\delta \frac{M_L}{M_X} M_R^2 \\ \lambda\delta \frac{M_L}{M_X} M_R^2 & \alpha M_1^2 - \delta M_R^2 \end{pmatrix} \tag{12}$$

Here αM_1^2 is a notation for a sum of terms coming from different terms in the Lagrangian.

According to the fine-tuning condition, one of matrix (12) eigenvalues must be of M_W order. In this case, the second $SU(2)_L$ -doublet will have a mass of M_R order. One can readily calculate the angle of η - ξ mixing:

$$\begin{aligned} \eta &= \cos\theta\, \eta' - \sin\theta\, \xi' \\ \xi &= \sin\theta\, \eta' + \cos\theta\, \xi' \\ tg\, 2\theta &= \lambda \frac{M_L}{M_X} \end{aligned} \tag{13}$$

If we assume that η' is the genuine Higgs field (this can always be attained by choosing signs of coupling constants in (12)), then $\langle \eta'^{*} \rangle \neq 0$, $\langle \xi'^{*} \rangle = 0$ and we'll obtain:

$$\langle \phi \rangle = \begin{pmatrix} \cos\theta \langle \eta' \rangle & 0 \\ 0 & -\sin\theta \langle \eta'^{*} \rangle \end{pmatrix}, \langle \tilde{\phi} \rangle = \begin{pmatrix} \sin\theta \langle \eta' \rangle & 0 \\ 0 & \cos\theta \langle \eta'^{*} \rangle \end{pmatrix} \tag{14}$$

With account of the fact that the second term in (6) is forbidden due to discrete symmetry, we come to the following expression for the mass ratio of b and t quarks:

$$\frac{m_b}{m_t} \approx \frac{1}{2} \lambda \frac{M_L}{M_X} \tag{15}$$

The estimate is correct to factors of tenth order. It follows from this estimate that M_c cannot differ very much from M_X , say, no more than by a factor of 10. In this case M_P which is between M_X and M_c cannot differ strongly from M_X either. No restrictions take place for $SU(2)_R$ -breaking scale M_R from the quark mass spectrum.

The obtained SO(10) breaking scheme, when the grand unification scale M_X, the left-right discrete symmetry breaking scale M_P and the quark--lepton symmetry breaking scale M_c are close to each other, gives for the proton lifetime results practically the same as those obtained in our previous work [2] . In [2] , the SO(10) breaking scheme looked as follows:

$SO(10) \xrightarrow{\langle(15,1,1)_{45}\rangle = M_X} G_1 \xrightarrow{\langle(\overline{10},1,3)_{126}\rangle = M_R} G_0$. The obtained renormalization group equations with account of the second loop enable us to calculate M_X and M_R versus the Weinberg angle. These results will be valid for the present model as well. With account of experimental value of the Weinberg angle, $\sin^2\theta_W = 0.22 \pm 0.02$, the least value of M_R (for the value of QCD Λ - parameter, 0.16 GeV) is $3\cdot 10^7$ GeV; to this corresponds $M_X = 10^{16}$ GeV and the proton lifetime, $\tau_p \leq 10^{37}$ yrs, which is much higher than experimental limit.

Note also that the SO(10) breaking scheme (10) can be generalized by considering the $SU(2)_R$ breaking in two steps. First, the $SU(2)_R$ is broken by v.e.v. $\langle(1, 1, 3)_{45}\rangle = M_R$ up to $U(1)_R$; then the v.e.v. $\langle(\overline{10}, 1, 3)_{126}\rangle = M_R^0$ breaks $U(1)_R \times U(1)_{B-L}$ up to $U(1)_y$. For such breaking scheme, the mixing of η , ξ doublets arises due to the same diagrams of Fig. 1. Only now, a nondiagonal in η , ξ mass term has the form: $\lambda\gamma \frac{M_c}{M_X} M_R^{02}(\eta^*\xi + \xi^*\eta)$. Finally, we'll obtain the same relation (14) for the mass ratio of b and t quarks. The difference is that the second $SU(2)_L$ doublet (except the Higgs field with a mass $\sim M_W$) arising at diagonalization of mass matrix η , ξ now will have a mass $\sim M_R^\circ$. Note that with the introduction of the new scale M_R° , the predictions for the proton lifetime change little, while small values of M_R° ($\sim$1 TeV) may bring to observable effects [4] .

Our consideration hitherto was restricted to the heaviest third generation of b , t quarks neglecting quark masses of the first two generations, u, d, s, c , and mixing. Now we want to include the first two generations and mixing effects. For that, we shall determine the transformation laws for three generations of fermions - the SO(10) $\underline{16}$-plets - with respect to our introduced discrete symmetry. If all the three generations transform identically:

$$\underline{16}_{1,2,3} \longrightarrow e^{-i\frac{\pi}{4}}\, \underline{16}_{1,2,3} \,, \tag{16}$$

then the interaction of fermions with the $\underline{10}$ Higgs field will have the form:

$$\sum_{i,j=1}^{3} \lambda_{ij}\, \underline{16}_i\, \underline{16}_j\, \underline{10} \;+\; \text{c.c.} \tag{17}$$

With respect to (14), we'll obtain the following relation between the mass matrices of up, u, c, t and down, d, s, b quarks: $\mathcal{M}_d = tg\,\theta\,\mathcal{M}_u$. Finally, we'll obtain that all mixing angles in the left charged current (Kobayashi-Maskawa angles) are zero and there take place equalities $\frac{m_d}{m_u} = \frac{m_s}{m_c} = \frac{m_b}{m_t} = tg\,\theta$. This is, of course, unacceptable. One can show that the situation cannot be improved by complication of Higgs sector either.

There exists, however, another possible way to determine fermion transformations relative to discrete symmetry, when not all generations of fermions transform identically:

$$\underline{16}_{1,3} \longrightarrow e^{-i\frac{\pi}{4}}\, \underline{16}_{1,3} \quad ; \qquad \underline{16}_2 \longrightarrow e^{i\frac{\pi}{4}}\, \underline{16}_2 \tag{18}$$

In this case, the interaction with the Higgs field $\underline{10}$ has the form:

$$\sum_{i,j=1}^{3} \left(\lambda_{ij}\, \underline{16}_i\, \underline{16}_j\, \underline{10} + \mu_{ij}\, \underline{16}_i\, \underline{16}_j \cdot \underline{10}^* \right) + c.c. \qquad (19)$$

where λ_{ij} and μ_{ij} are 3 x 3 matrices of the form:

$$\lambda = \begin{pmatrix} \lambda_1 & 0 & \lambda_4 \\ 0 & \lambda_2 & 0 \\ \lambda_4 & 0 & \lambda_3 \end{pmatrix}, \quad \mu = \begin{pmatrix} 0 & \mu_1 & 0 \\ \mu_1 & 0 & \mu_2 \\ 0 & \mu_2 & 0 \end{pmatrix} \qquad (20)$$

Then the mass matrices of the up and down quarks can be written as follows:

$$\mathcal{M}_u = (A + B\,tg\,\theta)\,e^{i\alpha}$$
$$\mathcal{M}_d = -(A\,tg\,\theta + B)\,e^{-i\alpha}$$

$$A = \begin{pmatrix} a_1 & 0 & a_4 \\ 0 & a_2 & 0 \\ a_4 & 0 & a_3 \end{pmatrix}, \quad B = \begin{pmatrix} 0 & b_1 & 0 \\ b_1 & 0 & b_2 \\ 0 & b_2 & 0 \end{pmatrix} \qquad (21)$$

Mass matrices of the (21) type were already considered **in models with** $SU(2)_L \times SU(2)_R$ symmetry [11] (see also [12] and references cited there). **After diagonalization** we'll obtain for the Kobayashi-Maskawa angles θ_1, θ_2, θ_3:

$$\frac{m_d}{m_s}\left(1 - \frac{m_u}{m_d}\frac{m_b}{m_t}\right) \leq \sin^2\theta_1 \leq \frac{m_d}{m_s}\left(1 + \frac{m_u}{m_d}\frac{m_b}{m_t}\right)$$
$$\left|\frac{m_c}{m_t} - \frac{m_s}{m_b}\right| \leq \sin^2\theta_2 \leq \frac{m_c}{m_t} + \frac{m_s}{m_b} \qquad (22)$$
$$\sin\theta_3 = \frac{m_s}{m_b}\sin\theta_2$$

this agreeing well **with experiment** [13] . For the parameter $tg\,\theta$ we obtain

$tg\vartheta \approx m_c/m_t$, in accordance with (13)-(15).

We can see that our introduced discrete symmetry (18) allows us to take account of mixing between existing three families of quarks and obtain a good prediction for the angles.

We can prove that the form of quark mass matrices (21), when fermions transform with respect to the action of discrete symmetry according to (18), is rather general: it does not change when fermions interact with several 10-plets or 126-plets.

Consider, e.g., the case when as Higgs fields we choose two representations $\underline{10}_1$ and $\underline{10}_2$ which transform similarly relative to discrete symmetry:

$$\underline{10}_1 \rightarrow i\,\underline{10}_1 \quad , \quad \underline{10}_2 \rightarrow i\,\underline{10}_2 \tag{23}$$

The interaction responsible for the fermion mass has the form:

$$\sum_{i,j=1}^{3} \big(\bar{Q}_{iL} \lambda^1_{ij} \Phi_1 Q_{jR} + \bar{Q}_{iL} \tilde{\Phi}_1 \mu^1_{ij} Q_{jR} + \bar{Q}_{iL} \lambda^2_{ij} \Phi_2 Q_{jR} + \bar{Q}_{iL} \tilde{\Phi}_2 \mu^2_{ij} Q_{jR} \big) + c.c. \tag{24}$$

$$Q_{iL,R} = \begin{pmatrix} u_i \\ d_i \end{pmatrix}_{L,R}$$

where λ^1, λ^2 and μ^1, μ^2 are matrices of the (20) type.

Consider the mass matrix of Φ_1, Φ_2 Higgs fields without account of terms of M_R^2 order. After its diagonalization we'll obtain new states Φ, Φ'. The interaction (24) is to be rewritten as:

$$\sum_{i,j=1}^{3} \big(\bar{Q}_{iL} \lambda_{ij} \Phi Q_{jR} + \bar{Q}_{iL} \mu_{ij} \tilde{\Phi} Q_{jR} + \bar{Q}_{iL} \lambda'_{ij} \Phi' Q_{jR} + \bar{Q}_{iL} \mu'_{ij} \tilde{\Phi}' Q_{jR} \big) + c.c. \tag{25}$$

where λ, λ', μ, μ' are some new matrices of the (20) type. According to the fine-tuning condition, one of the fields Φ, Φ', say, Φ, must have a mass $\ll M_X$ (i.e. $\sim M_R$). Evidently, precisely this field (up to terms M_R/M_X) contains the genuine Higgs. Then the third and fourth terms in (26) will not contribute to fermion masses, so we return back to the form of mass matrices contributed by (21).

Thus, we have proposed a model which, by means of additional discrete symmetry, enables one to avoid the contradiction with fermion mass spectrum arising in the SO(10) breaking schemes with intermediate $SU(4) \times SU(2)_L \times SU(2)_R$ or $SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$ symmetry. The obtained relation between b and t quark masses fixes a certain hierarchy of SO(10) breaking scales, the proton lifetime here predicted being in agreement with experiment. The introduced discrete symmetry at the same time allows one to obtain a phenomenologically acceptable model of quark mixing as well as to give good predictions for mixing angles in the left charged currents.

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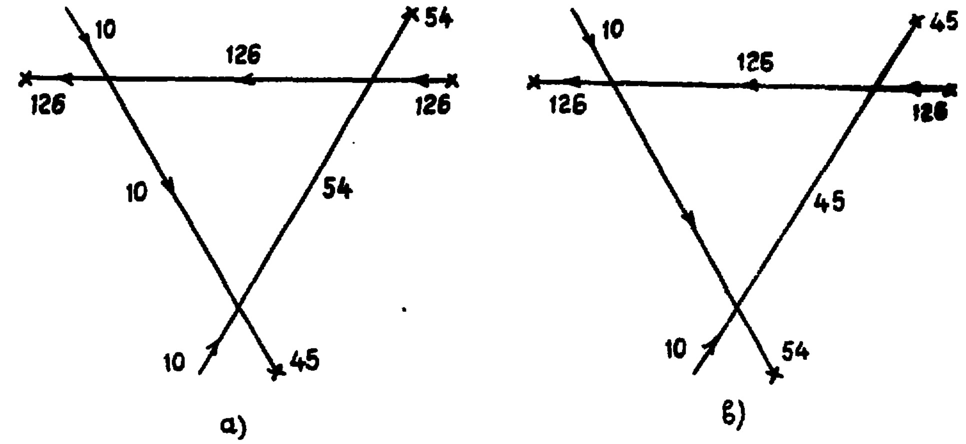

Fig.1. Diagrams contributing to a nondiagonal in η_L, ξ mass term in terms of fields - representations of SO(10).

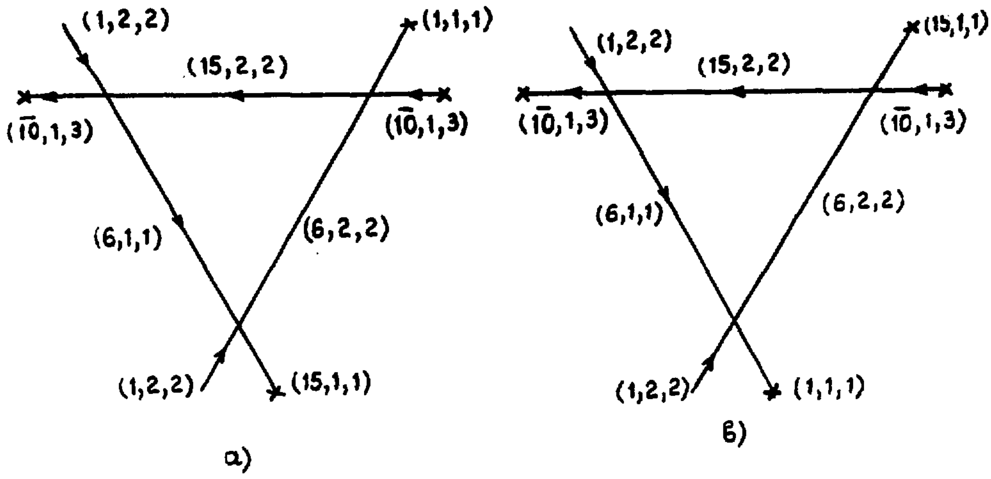

Fig.2. Diagrams contributing to a nondiagonal in η_L, ξ mass term in terms of fields - representations of $SU(4) \times SU(2)_L \times SU(2)_R$.

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МАССЫ КВАРКОВ

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