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RADIATIVE FERMION MASS MATRIX GENERATION IN SUPERSYMMETRIC MODELS

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ABSTRACT

Supersymmetric $SU(2)_L \otimes U(1)$ horizontal models are studied. The non-renormalization theorems of supersymmetry are used to make the mass generation and flavour mixing natural. For three families, we study the fermion mass matrix generation mechanism as a radiative effect due to horizontal interactions, using various representations of the gauge horizontal groups $SU(2)_H$ and $SU(3)_H$. An attractive possibility leading to a realistic mass matrix is found.

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A fundamental problem in the present theoretical description of elementary particle physics, as is expressed in the standard $SU(3)_C \otimes SU(2)_L \otimes U(1)$ picture¹⁾, is the plethora of free parameters. It is widely believed that there is a more symmetric world in higher energies and the remnant of its strong breaking is what we see in the presently accessible energies. Moreover, this high energy symmetric theory should contain very few parameters, while all the undetermined parameters of the low energy standard model should be calculable. The grand unification schemes²⁾ are well-known attempts in this direction; however, they suffer from various diseases, the most serious of which is the hierarchy problem³⁾.

In the present work, our aim is to obtain some insight into the fermion mass matrix, and we speculate that the relevant energy scale is of a few hundred TeV. Thus, we demand that the theory has also, in addition to the usual $SU(3)_C \otimes SU(2)_L \otimes U(1)$ gauge symmetry, horizontal gauge symmetry and global supersymmetry at that energy scale.

Our approach in using horizontal gauge interactions^{4),5)} is to give a direct mass only to the heavy third generation through Yukawa couplings, and then to produce the masses of the lighter ones and the mixing angles radiatively⁵⁾ by exchanges of heavy horizontal vector superfields. If such a programme is to have any hope of helping the construction of realistic models of weak and electromagnetic interactions, it must first successfully confront two well-known interrelated problems: naturalness in the mass generation mechanism, and no fine tuning in the Higgs sector of the theory.

The principle of naturalness can be stated as follows⁶⁾: a theory is natural if it does not contain small quantities. However, even if a theory contains some small quantities, it can still be natural if, for every quantity which almost vanishes, a symmetry exists which, if unbroken, would require that quantity to vanish exactly. This concept of naturalness, when it is applied to the fermion mass generation in renormalizable gauge theories with spontaneous symmetry breaking, is realized by constraining the Yukawa sector of the theory by some symmetries called natural symmetries⁷⁾. Various models^{7),8),5)}, more or less successful, which incorporate such natural symmetries, appear in the literature.

The problem of the fine tuning here comes about from the requirement of having a fermion mass matrix at tree level of the form (for three families):

$$M = m \cdot \text{diag}(0, 0, 1) \quad (1)$$

Then, as we shall discuss subsequently, an adjustment of parameters is generally necessary.

Supersymmetry⁹⁾, because of its good behaviour under renormalization¹⁰⁾, has been recently introduced in GUTs¹¹⁾ in order to help find the solution of the hierarchy problem. The concept of naturalness in supersymmetric theories is somehow different than in non-supersymmetric gauge theories¹²⁾, but what concerns us here is that the non-renormalization theorems¹⁰⁾ tell us that in supersymmetric theories, one can impose arbitrary relations among mass and Yukawa coupling parameters, or set them equal to zero provided that supersymmetry is unbroken. The reason is that only a universal wave function renormalization for all the parameters in the theory is generated in perturbation theory, and this is also true in the presence of an internal gauge symmetry¹³⁾.

In this paper, we propose that supersymmetry (in addition to gauge symmetries) protects the appearance of various matrix elements and its breaking produces finite matrix elements proportional to it.

Recently it was pointed out¹⁴⁾ that in supersymmetric GUTs it is not possible to generate realistic fermion masses radiatively. We shall show that one can obtain radiative fermion masses in a supersymmetric horizontal model when supersymmetry and horizontal symmetry break spontaneously at intermediate mass scales^{*)}. Consider that a model, apart from matter superfields Q and vector superfields V , contains a set of light chiral superfields ϕ , which carry $SU(2)_L \otimes U(1)$ and horizontal quantum numbers and play the rôle of the Higgs fields which break $SU(2)_L \otimes U(1)$ down to $U(1)$ and generate mass for the third heavy family. Assume there also exists a set of heavy chiral superfields Σ with only horizontal quantum numbers, which break the horizontal symmetry, and that supersymmetry breaking is realized by an F type breaking mechanism¹⁶⁾ with a set of singlet superfields X .

The scalar bosons of such a model will obtain radiative masses through diagrams of the form shown in Fig. 1. One finds

$$m_{A_q}^2 \sim g^n \frac{\mu^4}{\mu_H^2} Q^* Q \quad (2)$$

^{*)} A similar observation has been made in Ref. 15).

where $\mu^2 = F_X$, g^n stands for the coupling constants and μ_H is the heavy horizontal mass^{*)} which circulates inside the diagram. If $\mu^2 \sim \mu_W \mu_H$, then $m_{Aq} \sim \mu_W$, which is the desired order of magnitude for the Higgs masses.

Quarks and leptons will obtain radiative masses through the diagram of Fig. 2. One finds

$$m_{\psi q} \sim \alpha_H f \frac{\langle F_\varphi \rangle}{\mu_H} \quad (3)$$

where α_H is the horizontal fine coupling constant and f is the Yukawa coupling of the third generation. In order to obtain a non-zero contribution from the diagram of Fig. 2, we must have $\langle F_\varphi \rangle \neq 0$, while in order to obtain a realistic contribution, we must have $\langle F_\varphi \rangle \approx \langle A_\varphi \rangle \langle A_\Sigma \rangle$. In this case, Eq. (3) gives

$$m_{\psi q} \sim \alpha_H \cdot m \quad (4)$$

where m is the mass of the heaviest fermion of the same charge. As we shall discuss in the following, this condition is a very strong constraint for model building, and it can only be satisfied if in the scalar potential there is a direct coupling of φ with Σ .

In our effort to materialize the above picture, we have studied supersymmetric models based on the gauge group $SU(2)_L \otimes U(1) \otimes G_H$, with $G_H = SU(2)_H, SU(3)_H$. Our strategy is to produce in a natural way a mass matrix for fermions of the same charge of the form of Eq. (1) and then, when the theory is spontaneously broken, the radiative corrections add such contributions that the mass matrix becomes

$$M = m \begin{pmatrix} \alpha_H^2 & \alpha_H^2 & \alpha_H^2 \\ \alpha_H^2 & \alpha_H & \alpha_H^2 \\ \alpha_H^2 & \alpha_H^2 & 1 \end{pmatrix} \quad (5)$$

which is a good approximation of the observed fermion mass pattern.

^{*)} From the experimental bounds on the flavour-changing neutral currents $\mu_H \sim 10-100 \text{ TeV}$ ^{4), 5)}.

Let us start with $G_H = SU(2)_H$. The superfields involved and their representations are given in Table I. The superpotential of the model is given by

$$\begin{aligned} W = & M(S^T \epsilon \bar{H} + H^T \epsilon \bar{S}) + \lambda Y(S^T \epsilon \bar{S} - \mu^2) \\ & + k_1 Y_1(\text{Tr} \Sigma \bar{\Sigma} - \pi_1^2) + k_2 Y_2(X \bar{X} - \pi_2^2) \\ & + f(Q^{T\alpha} \epsilon X D^\alpha) + f'(Q^{T\alpha} \epsilon \bar{X} U^\alpha) \\ & + g(Q^{T\alpha} D^\alpha) \epsilon \Sigma_{\alpha\beta} + g'(Q^{T\alpha} U^\alpha) \epsilon \bar{\Sigma}_{\alpha\beta} \end{aligned} \quad (6)$$

where $\epsilon = i\sigma^2$; $\alpha, \beta = 1, 2, 3$ family indices.

From the very beginning, this model suffers from two serious problems. Assume that the scalar components of $X(\bar{X})$ and $\Sigma(\bar{\Sigma})$ develop the following vacuum expectation values (vev's):

$$\langle A_\Sigma \rangle = \langle A_{\bar{\Sigma}} \rangle = \begin{pmatrix} V & & \\ & V & \\ & & -2V \end{pmatrix}; \quad \langle A_X \rangle = \langle A_{\bar{X}} \rangle = V \quad (7)$$

The fermion mass matrix acquires non-zero contributions from the vev's in Eq. (7) and in order to obtain the form given in Eq. (1), one has to demand

$$fV + gV = 0 \quad (8)$$

The same result would have held if we had used triplets and quintets for $X(\bar{X})$ and $\Sigma(\bar{\Sigma})$ respectively under the $SU(2)_H$. Thus the fine tuning is unavoidable in the $SU(2)_H$ model. The more serious problem comes about when one minimizes the potential

$$V = \sum_A \left| \frac{\partial W}{\partial A} \right|^2 + \frac{1}{2} \sum_K |D^K|^2 \quad (9)$$

where A are the scalar components of the chiral superfields and

$D_K = \sum_{\alpha, \beta} (t_K)_{\alpha\beta}^{A^*} A^\alpha A^\beta$, with t_K the generators of the gauge groups and the D term^{*)}. Then we found that supersymmetry and horizontal symmetry are spontaneously broken,

^{*)} In this model, because the representations appear in pairs of opposite chirality and develop equal vev's, they give vanishing contributions to the D term.

but $\langle F_{X,\bar{X},\Sigma,\bar{\Sigma}} \rangle = 0$. This is again true for $X,\bar{X}; \Sigma,\bar{\Sigma}$ being triplets and quintets under $SU(2)_H$. The reason we obtain this result is that gauge invariance forbids the appearance of couplings between $X,\bar{X},\Sigma,\bar{\Sigma}$ and $S,\bar{S},H,\bar{H}$, and only if the latter superfields belong to a high representation under $SU(2)_H$ would it be possible, in principle, to avoid such a situation.

Next let us consider a supersymmetric model with $G_H = SU(3)_H$. The superfield content is given in Table II^{*)}. The superpotential of the model (with omission of the Yukawa terms) is

$$\begin{aligned} W = & m_\varphi \varphi'_{\alpha\beta} \varphi^{\alpha\beta} + g_\varphi \varphi'_{\alpha\beta} \Sigma^\beta_\gamma \varphi^{\gamma\alpha} + m_\Sigma \Sigma^\alpha_\beta \Sigma^\beta_\alpha + g_\Sigma \Sigma^\alpha_\beta \Sigma^\beta_\gamma \Sigma^\gamma_\alpha \\ & + m_P \Sigma^\alpha_\beta P^\beta_\alpha + m_H H'_\alpha H^\alpha + m_T T'_\alpha T^\alpha + g_H H'_\alpha \Sigma^\alpha_\beta H^\beta \\ & + g_T T'_\alpha \Sigma^\alpha_\beta T^\beta + \lambda Y (\Sigma^\alpha_\beta \Sigma^\beta_\alpha - \mu^2) + \omega X (\varphi'_{\alpha\beta} \varphi^{\alpha\beta} - v^2) \end{aligned} \quad (10)$$

This superpotential leads to a spontaneous breaking of the horizontal symmetry and global supersymmetry. Here we shall not present the potential and its full minimization, but we rather concentrate on the coupling of the scalar components of φ and Σ superfields as is expressed through the minimization equations

$$(m_\varphi + \omega A_X) A_\varphi^{\alpha\beta} + \frac{1}{2} g_\varphi (A_{\Sigma\gamma}^\alpha A_\varphi^{\gamma\beta} + A_{\Sigma\delta}^\beta A_\varphi^{\delta\alpha}) = 0 \quad (11)$$

We can always use an $SU(3)$ rotation to bring $A_\varphi^{\alpha\beta}$ in a diagonal form

$$\langle A_\varphi^{\alpha\beta} \rangle = v_\alpha \delta^{\alpha\beta} \quad (12)$$

Then, from Eq. (11), we obtain

$$\langle A_{\Sigma\beta}^\alpha \rangle = a_\alpha \delta^\alpha_\beta \quad (13)$$

^{*)} The introduction of the sextet representation under $SU(3)_H$ leads to anomalies in the $SU(3)_H$ currents. These anomalies can be cancelled by the introduction of heavy fermions, which obtain masses at the scale of horizontal symmetry breaking.

Using Eqs. (12) and (13), Eq. (11) becomes

$$\sum_{i=1}^3 \left(m_\varphi + \omega \langle A_x \rangle + g_\varphi a_i \right) v_i = 0 \quad (14)$$

If now $v_i \neq 0$ for every i , then $a_1 = a_2 = a_3$, in contradiction to the traceless condition of the A_Σ . Thus, we can always choose

$$\langle A_\varphi^{\alpha\alpha} \rangle = \text{diag}(0, 0, v) \quad (15)$$

with

$$a_3 = - \frac{m_\varphi + \omega \langle A_x \rangle}{g_\varphi} \quad (16)$$

Equation (15) says that only the third family will obtain mass at tree level. If we allow the scalar components of P, T, T', H, H' superfields to develop vev's in the direction

$$\begin{aligned} \langle A_P^{\alpha\alpha} \rangle &= P_\alpha \delta^{\alpha\alpha} \\ \langle A_H^{\alpha\alpha} \rangle &= \langle A_{H\alpha}^{\prime\alpha} \rangle = n \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix} \\ \langle A_T^{\alpha\alpha} \rangle &= \langle A_{T\alpha}^{\prime\alpha} \rangle = \tau \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix} \end{aligned} \quad (17)$$

we find $\langle F_\varphi \rangle = \langle F_{\varphi'} \rangle = v \cdot a$, but unfortunately the diagram in Fig. 2, as we find by analyzing it in superfield components, gives equal contributions to $d(u)$ and $s(c)$ quarks

$$m_{d(u)} = m_{s(c)} \sim \alpha_H m_{b(t)} \quad (18)$$

Note that if we were to use octets and singlets instead of octets and sextets, we would face the fine tuning problem at the tree level, as in the $SU(2)_H$ case.

In the following, we shall present a supersymmetric $SU(2)_L \otimes U(1) \otimes SU(3)_H$ model which overcomes most of the difficulties we faced in the previous attempts. The list of the superfield contents of this model is given in Table III.

We consider the following superpotential:

$$\begin{aligned}
 W = & \lambda_X (\text{Tr } \Sigma^2 - \mu^2) + \lambda_T \text{Tr } \Sigma^2 T \\
 & + m_\varphi \text{Tr } \varphi' \varphi + g_\varphi \text{Tr } \varphi' \Sigma \varphi + g_\varphi' \text{Tr } \varphi' \Sigma' \varphi \\
 & + m_6 \text{Tr } \Sigma^2 + g_6 \text{Tr } \Sigma^3 + m_6' \text{Tr } \Sigma'^2 + g_6' \text{Tr } \Sigma'^3 \\
 & + m \text{Tr } \Sigma \Sigma' + \rho_6 \text{Tr } \Sigma^2 \Sigma' + \rho_6' \text{Tr } \Sigma'^2 \Sigma \\
 & + m_H H H' + g_H H' \Sigma H + g_H' H' \Sigma' H \\
 & + K_1 Y (\text{Tr } \varphi \varphi' - v^2) + K_2 Q Q^u \varphi + K_3 Q Q^d \varphi'
 \end{aligned} \tag{19}$$

The scalar potential is given in Eq. (8). Again, we are not going to present the potential^{*)} and minimization conditions, but we will rather state the results of our examination. We find that global supersymmetry breaks down spontaneously and the scalar components of the chiral superfields can develop the following vev's:

$$\langle A_\varphi \rangle = \frac{v}{2\sqrt{6}} (\lambda_6 + i\lambda_7) \quad ; \quad \langle A_{\varphi'} \rangle = \frac{v}{2\sqrt{6}} (\lambda_6 - i\lambda_7) \tag{20}$$

$$\left. \begin{aligned}
 \langle A_\Sigma \rangle = \langle A_T \rangle &= \frac{V}{2} \lambda_6 \quad ; \quad \langle A_{\Sigma'} \rangle = \frac{V'}{2} \lambda_7 \\
 \langle A_H \rangle = \langle A_{H'} \rangle &= n \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}
 \end{aligned} \right\} \tag{21}$$

$$\langle A_q \rangle = \langle A_{q^u} \rangle = \langle A_{q^d} \rangle = 0 \tag{22}$$

with $V, V', n \gg v$

which in turn completely break the $SU(3)_H$ and $SU(2)_L \otimes U(1)$ to $U(1)$ electromagnetic.

^{*)} Because of the reality of the eight-dimensional adjoint representations of $SU(3)_H$, their D terms are zero¹⁷⁾. For the rest of the representations, see the third footnote.

The quark mass matrix, after symmetry breaking, takes at tree level the form:

$$\left(\bar{q}_{01}^d, \bar{q}_{02}^d, \bar{q}_{03}^d \right)_L \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & m \\ 0 & (m') & 0 \end{pmatrix} \begin{pmatrix} q_{01}^d(u) \\ q_{02}^d(u) \\ q_{03}^d(u) \end{pmatrix}_R + h.c. \quad (23)$$

where $m = K_3/2\sqrt{6} v$ and $m' = K_2/2\sqrt{6} v$, q stands for quarks and the subscript zero indicates that the q 's are not yet the physical ones. In this order in perturbation theory, we may diagonalize the above mass matrix

$$\begin{aligned} & \bar{D}_{0L}(\bar{u}_{0L}) \mathcal{M}_{D(u)} D_{0R}(u_{0R}) \\ &= \bar{D}_{0L}(\bar{u}_{0L}) (S^{-1} S) \mathcal{M}_{D(u)} S^{-1} S D_{0R}(u_{0R}) \end{aligned} \quad (24)$$

where

$$S = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix} \quad (25)$$

Diagonalization leads us to make the following identifications in this order in perturbation theory:

$$\begin{aligned} (q_{01}^d, q_{02}^d, q_{03}^d)_L &\equiv (d_0, s_0, b_0)_L; (q_{01}^d, q_{02}^d, q_{03}^d)_R \equiv (d_0, b_0, s_0)_R \\ (q_{01}^u, q_{02}^u, q_{03}^u)_L &\equiv (u_0, t_0, c_0)_L; (q_{01}^u, q_{02}^u, q_{03}^u)_R \equiv (u_0, c_0, t_0)_R \end{aligned} \quad (26)$$

and b and t quarks obtain masses m and m' respectively at tree order in a natural way without any fine tuning.

Next, looking at the one-level contributions to the mass matrix, we find that they are realistic, since $\langle F_\phi \rangle \sim v \cdot V$, $\langle F_{\phi'} \rangle \sim v \cdot V'$ according to our discussion that led us to Eq. (4). Moreover, because left- and right-handed quarks appear asymmetrically⁵⁾ [see Eq. (26)], we find that only the second family receives

radiative masses at this order

$$m_{S(c)} \sim \alpha_H \cdot m(m') \quad (27)$$

In two loops, of course, we expect the various matrix elements of the quark matrix to obtain contributions from diagrams such as in Fig. 3. Thus, the mass matrix for quarks of the same charge takes the form of Eq. (5). However, because of the plethora of parameters that are involved in the model, it is very difficult to check that the vev's given in Eqs. (20)-(22) correspond to the deepest possible vacuum.

In conclusion, we have presented a study of supersymmetric $SU(2)_L \otimes U(1) \otimes G_H$ models, in which the various problems related to radiative mass generation were analyzed. We found an attractive supersymmetric $SU(2)_L \otimes U(1) \otimes SU(3)_H$ model in which it is possible to generate fermion masses and mixing angles radiatively in a natural way.

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TABLE I

<u>Superfield</u>	<u>Components</u>	<u>$SU(2)_L \otimes U(1) \otimes SU(2)_H$</u> <u>Quantum numbers</u>	
a) <u>Vector superfields</u>			
V_W	$(V_W; \lambda_W; D_W)$	$(3, 0, 1)$	$SU(2)_L$ vector bosons
V_Y	$(V_Y; \lambda_Y; D_Y)$	$(1, 1, 1)$	$U(1)$ vector boson
V_H	$(V_H; \lambda_H; D_H)$	$(1, 0, 3)$	$SU(2)_H$ horizontal vector bosons
b) <u>Fermion chiral superfields</u>			
Q	(A_q, ψ_q, F_q)	$(2, 1/3, 3)$	left-handed quark doublets
U	$(A_{qu}, \psi_{qu}, F_{qu})$	$(1, 4/3, 3)$	right-handed up- quarks
D	$(A_{qd}, \psi_{qd}, F_{qd})$	$(1, -2/3, 3)$	right-handed down- quarks
c) <u>Higgs chiral superfields</u>			
S	(A_S, ψ_S, F_S)	$(1, 0, 2)$	Higgs fields respon- sible for $SU(2)_H$ breaking
$\bar{S}$	$(A_{\bar{S}}, \psi_{\bar{S}}, F_{\bar{S}})$	$(1, 0, 2)$	
H	(A_H, ψ_H, F_H)	$(1, 0, 2)$	
$\bar{H}$	$(A_{\bar{H}}, \psi_{\bar{H}}, F_{\bar{H}})$	$(1, 0, 2)$	
X	(A_X, ψ_X, F_X)	$(2, 1, 1)$	Higgs fields responsible for $SU(2)_L \otimes U(1)$ breaking and for the mass of the heaviest fermion
$\bar{X}$	$(A_{\bar{X}}, \psi_{\bar{X}}, F_{\bar{X}})$	$(2, -1, 1)$	
Σ	$(A_{\Sigma}, \psi_{\Sigma}, F_{\Sigma})$	$(2, 1, 5)$	
$\bar{\Sigma}$	$(A_{\bar{\Sigma}}, \psi_{\bar{\Sigma}}, F_{\bar{\Sigma}})$	$(2, -1, 5)$	
Y	(A_Y, ψ_Y, F_Y)	$(1, 0, 1)$	Singlets
Y_1	$(A_{Y_1}, \psi_{Y_1}, F_{Y_1})$	$(1, 0, 1)$	
Y_2	$(A_{Y_2}, \psi_{Y_2}, F_{Y_2})$	$(1, 0, 1)$	

TABLE II

<u>Superfield</u>	<u>Components</u>	<u>$SU(2)_L \otimes U(1) \otimes SU(3)_H$</u> <u>Quantum Numbers</u>	
a) <u>Vector superfields</u>			
V_W	(V_W, λ_W, D_W)	$(3, 0, 1)$	$SU(2)_L$ vector bosons
V_Y	(V_Y, λ_Y, D_Y)	$(1, 1, 1)$	$U(1)$ vector boson
V_H	(V_H, λ_H, D_H)	$(1, 0, 8)$	$SU(3)_H$ horizontal Vector bosons
b) <u>Fermion chiral superfields</u>			
Q	(A_q, ψ_q, F_q)	$(2, 1/3, 3)$	left-handed quark doublet
U	$(A_{qu}, \psi_{qu}, F_{qu})$	$(1, 4/3, \bar{3})$	right-handed up- quarks
D	$(A_{qd}, \psi_{qd}, F_{qd})$	$(1, -2/3, \bar{3})$	right-handed down- quarks
c) <u>Higgs chiral superfields</u>			
Σ	$(A_\Sigma, \psi_\Sigma, F_\Sigma)$	$(1, 0, 8)$	Higgs fields respon- sible for $SU(3)_H$ breaking
P	(A_P, ψ_P, F_P)	$(1, 0, 8)$	
T	(A_T, ψ_T, F_T)	$(1, 0, 3)$	
T'	$(A_{T'}, \psi_{T'}, F_{T'})$	$(1, 0, \bar{3})$	
H	(A_H, ψ_H, F_H)	$(1, 0, 3)$	
H'	$(A_{H'}, \psi_{H'}, F_{H'})$	$(1, 0, \bar{3})$	
φ	$(A_\varphi, \psi_\varphi, F_\varphi)$	$(2, 1, 6)$	Higgs responsible for $SU(2)_L \otimes U(1)$ break- ing and for the mass of heaviest fermions.
φ'	$(A_{\varphi'}, \psi_{\varphi'}, F_{\varphi'})$	$(2, -1, \bar{6})$	
Y	(A_Y, ψ_Y, F_Y)	$(1, 0, 1)$	Singlets
X	(A_X, ψ_X, F_X)	$(1, 0, 1)$	

TABLE III

<u>Superfield</u>	<u>Components</u>	<u>$SU(2)_L \otimes U(1) \otimes SU(3)_H$</u> <u>Quantum Numbers</u>	
a) <u>Vector superfields</u>			
V_W	$(V_W; \lambda_W; D_W)$	$(3, 0, 1)$	$SU(2)_L$ vector bosons
V_Y	$(V_Y; \lambda_Y; D_Y)$	$(1, 1, 1)$	$U(1)$ vector boson
V_H	$(V_H; \lambda_H; D_H)$	$(1, 0, 8)$	$SU(3)_H$ horizontal vector bosons
b) <u>Fermion chiral superfields</u>			
Q	(A_q, ψ_q, F_q)	$(2, 1/3, 3)$	left-handed quark doublets
U	$(A_{qu}, \psi_{qu}, F_{qu})$	$(1, 4/3, 3)$	right-handed up- quarks
D	$(A_{qd}, \psi_{qd}, F_{qd})$	$(1, -2/3, 3)$	right-handed down- quarks
c) <u>Higgs chiral superfields</u>			
Σ	$(A_\Sigma, \psi_\Sigma, F_\Sigma)$	$(1, 0, 8)$	} Higgs responsible for the breaking of the horizontal gauge group
Σ'	$(A_{\Sigma'}, \psi_{\Sigma'}, F_{\Sigma'})$	$(1, 0, 8)$	
Σ''	$(A_{\Sigma''}, \psi_{\Sigma''}, F_{\Sigma''})$	$(1, 0, 8)$	
H	(A_H, ψ_H, F_H)	$(1, 0, 3)$	
H'	$(A_{H'}, \psi_{H'}, F_{H'})$	$(1, 0, \bar{3})$	
φ	$(A_\varphi, \psi_\varphi, F_\varphi)$	$(2, 1, 8)$	} Higgs responsible for $SU(2)_L \otimes U(1)$ breaking and for the mass of heaviest fermions
φ	$(A_{\varphi'}, \psi_{\varphi'}, F_{\varphi'})$	$(2, 1, 8)$	
X	(A_X, ψ_X, F_X)	$(1, 0, 1)$	} Singlets
Y	(A_Y, ψ_Y, F_Y)	$(1, 0, 1)$	

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FIGURE CAPTIONS

- Fig. 1 : Diagram contributing to scalar masses.
- Fig. 2 : Diagram contributing to fermion masses.
- Fig. 3 : Two-loop diagram contributing to light family fermion masses and off-diagonal matrix elements.

FIG. 1

FIG. 2

FIG. 3