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Light↔Heavy Symmetry:

Geometric Mass Hierarchy for Three Families

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Abstract

The Universal Seesaw pattern coupled with a Light↔Heavy symmetry principle leads to the Diophantine equation $N = \sum_{i=1}^{N} n_i$, where $n_i \geq 0$ and distinct. Its unique non-trivial solution $(3 = 0+1+2)$ gives rise to the geometric mass hierarchy m_W, $m_W\epsilon$, $m_W\epsilon^2$ for $N = 3$ fermion families. This is realized in a model where the hybrid (yet Up↔Down symmetric) quark mass relations $m_d m_t \approx m_c^2 \leftrightarrow m_u m_b \approx m_s^2$ play a crucial role in expressing the CKM mixings in terms of simple mass ratios, notably $\sin\theta_C \approx \frac{m_c}{m_b}$.

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The physics which governs the Yukawa sector is rooted beyond the standard $SU(3) * SU(2) * U(1)$ electro/weak theory. Despite the fact that quark (lepton) masses and mixings are fairly well known, their observed mass hierarchy does not have at the moment a solid theoretical ground, not even a reasonable empirical formulation. The finest ideas around have so far fallen short of decoding the three family Fermi Puzzle. Supersymmetry (R-symmetry included) does not allow the same scalar couple in both Up and Down sectors, but otherwise leaves the Yukawa couplings fully arbitrary. Grand unification (GUT), attainable at the single-family level, has consequently little to say about the horizontal features of the fermion spectrum. And superstring theory, loaded with self-consistency and inspiration, has yet to become the theory of everything (TOE). In fact, it appears as though one must patiently wait the entry of (quantum) gravity into the game. In the meantime, it always makes sense to pay serious attention to the accumulating experimental clues.

In this letter, we introduce the so-called Light-Heavy symmetry in an attempt to account for the observed geometric Fermi mass hierarchy. The natural framework to host such a symmetry principle is the Universal Seesaw (US) model [1]. In this unifiable model, a Froggatt-Nielsen-type [2] mechanism was implemented in a universal manner, without appealing to any hierarchy among the non-vanishing Yukawa couplings. It was originally designed, using a simplified 'square root' Higgs system, to actually predict the Gell-Mann-Yanagida [3] $m_\nu \ll m_e$ once having accounted for $m_{e,u,d} \ll m_W$. In what follows, we shall show that the pattern of the US mechanism combined with the Light-Heavy symmetry idea points uniquely to $N = 3$ fermion families, and dictates a geometric hierarchy among their masses. This is achieved without upsetting the symmetric interplay between the Up and the Down sectors. We present the arguments in three steps:

<u>Step I:</u> Consider a typical US mass sub-matrix of the form

$$\begin{pmatrix} & & & & & m \\ & & & & \chi & M \\ & & & . & M & \\ & & . & . & & \\ & \chi & . & & & \\ \chi & M & & & & \end{pmatrix} , \tag{1}$$

where m denotes the electro/weak mass scale. χ and M are the two $SU(3) * SU(2) * U(1)$-invariant heavy mass scales, $m \ll \chi \ll M$, whose ratio defines the US hierarchy parameter $\epsilon \equiv \frac{\chi}{M}$. If n seesaw partners (weak singlets with matching $SU(3)_c * U(1)_{e.m}$ assignments) are involved, the lightest eigenvalue of the above $(n+1)$-dimensional sub-matrix is of order $m\epsilon^n$.

Step II: Requiring an arbitrary mass hierarchy among N standard families, to be precise, $m\epsilon^{n_1}, m\epsilon^{n_2}, ..., m\epsilon^{n_N}$ such that $n_i \neq n_j$ if $i \neq j$, simply means introducing a total number $n_1 + n_2 + ... + n_N$ of exotic seesaw families into the theory. Further, if we impose a Light-Heavy symmetry principle to pair one seesaw partner F with each standard fermion f, we obtain the Diophantine equation

$$n_1 + n_2 + ... + n_N = N \ , \tag{2}$$

to be satisfied by a set $n_i (i = 1, ..., N)$ of distinct non-negative integers.

Step III: The one-family solution $(N = 1; n_1 = 1)$ constitutes the original US model. This solution as well as the two-family solution $(N = 2; n_1 = 0, n_2 = 2)$ are nothing but the building blocks of the only non-trivial solution $(N = 3; n_1 = 0, n_2 = 1, n_3 = 2)$. Thus, not only have we correlated the total number $N = 3$ of families with the Fermi mass hierarchy, but we can also infer that

(i) Owing to $n_1 = 0$, the heaviest standard family necessarily picks up the electro/weak mass scale (an encouraging result given the top mass $m_t \approx 2m_W$), and

(ii) Owing to $n_1 + n_3 = 2n_2$, the hierarchy is necessarily geometric.

We now proceed to construct explicitly the three-family model. Our 6x6 mass texture consists of the three blocks associated with $n_1 = 0, n_2 = 1, n_3 = 2$ of the prototype form (1). But how are the remainder entries (to be referred to as block mixings) to be decided? After all, they can have undesirable (as well as desirable) consequences in the low-energy regime. A first reasonable criterion would be that the new entries should not introduce leading order corrections to the hierarchical eigenvalues $m_W, m_W\epsilon, m_W\epsilon^2$. A second consideration would be whether the resulting matrix has any symmetries left and whether it produces the right order of mixings. To illustrate these points, let us consider the two-family 3x3 sub-texture

$$\left(\begin{array}{cc|c} 0 & \tilde{m} & m \\ \hline 0 & m & 0 \\ \chi & M & \tilde{\chi} \end{array}\right) , \tag{3}$$

where the carefully located zeroes assure that in the limit $M \to \infty$ there is no vestige of the second family. Next, if both block mixings $\tilde{m}$ and $\tilde{\chi}$ are non-vanishing, there is no possibility for a remnant symmetry that predicts this pattern. Consequently, at least one of them should be zero. If $\tilde{\chi} = 0$, we can show that the mixing angle θ_L between the two light left-handed fermions is $\theta_L \sim \mathcal{O}(\epsilon^2)$. If, on the other hand, $\tilde{m} = 0$, we derive $\theta_L \sim \mathcal{O}(\epsilon)$. The latter choice is preferred if one wants to avoid the situation where the nearest neighbor mixings get doubly-suppressed. One need not be discouraged though by the fact that, unlike in the conventional Weinberg-Fritzsch prescription [4] for relatively large (square-mass-ratio) mixings $\sim \sqrt{\frac{\lambda_i}{\lambda_j}}$ (given $\lambda_i \ll \lambda_j$), our approach can only offer apparently small (mass-ratio) mixings $\sim \frac{\lambda_i}{\lambda_j}$.

Using similar arguments, we are led to a unique 6x6 mixing pattern for three standard families, and would like to support it by a specific Light-Heavy symmetry realization. We aim towards the Up-Down symmetric spectrum

$$\begin{aligned} &m_t \approx m &, \quad &m_b \approx n &, \\ &m_c \approx m \mid \tfrac{x}{M} \mid &, \quad &m_s \approx n \mid \tfrac{y}{M} \mid &, \\ &m_u \approx n \mid \tfrac{y}{M} \mid^2 &, \quad &m_d \approx m \mid \tfrac{x}{M} \mid^2 &, \end{aligned} \tag{4}$$

supposed to hold at some common (yet to be specified) mass scale. The reasons for considering such a hybrid geometric mass hierarchy are threefold: (i) The above alternative seems to be numerically preferred, (ii) Since $m_e m_\tau \approx m_\mu^2$ relation is badly violated in the charged lepton sector, there is no reason for insisting on $m_u m_t \approx m_c^2 \leftrightarrow m_d m_b \approx m_s^2$ in the quark sector, and (iii) $\frac{m}{n}$ may serve to enhance the Cabibbo angle (that is $\theta_c \sim \frac{m}{n}\frac{x}{M}$, rather than $\theta_c \sim \frac{x}{M}, \frac{y}{M}$). Assuming that all non-vanishing Yukawa couplings are of the same order of magnitude, $m \equiv < \phi_1 >$ and $n \equiv < \phi_2 >$ are essentially the VEVs of the two Higgs doublets involved, whereas $x \equiv < \chi_1 >$ and $y \equiv < \chi_2 >$ are the VEVs of their singlet companions, respectively. Recall that doublets and singlet Higgs scalars play a perfectly symmetric Yukawa role [1] in the US mechanism.

To derive the above pattern, we appeal to a horizontal $U(1)_Q$ global symmetry (soon to be discretized on electro/weak grounds). A central role in our analysis is played by the 6x6 matrix

$$Q_{ij}^{(\Psi)} \equiv Q(\Psi_{iL}) - Q(\Psi_{iR}) \ , \tag{5}$$

defined for each electrically-charged fermion sector. Denoting the Q-charges of the scalars by $Q(\phi_1) \equiv \alpha$, $Q(\phi_2) \equiv \beta$, $Q(\chi_1) \equiv a$, $Q(\chi_2) \equiv b$, the rules of the game are quite simple:

$$\begin{array}{lll} \text{for} \quad \Psi_L = u_L : & Q_{ij}^{up} = \alpha, \beta \Rightarrow M_{ij}^{up} = m, n & , \\ \text{for} \quad \Psi_L = d_L : & Q_{ij}^{down} = -\beta, -\alpha \Rightarrow M_{ij}^{down} = n^*, m^* & , \\ \text{for} \quad \Psi_L = F_L : & Q_{ij} = a, b, -a, -b \Rightarrow M_{ij} = x, y, x^*, y^* & . \end{array} \tag{6}$$

Notice the fine differences of the $SU(3)*SU(2)*U(1)$ restrictions on the Higgs singlet versus doublet couplings. For example, $\overline{U_L}\chi u_R$ and $\overline{U_L}\chi^\dagger u_R$ are both allowed. However, in the case of ordinary quarks, $\overline{q_L}\phi u_R$ and $\overline{q_L}\phi^\dagger d_R$ are allowed, $\overline{q_L}\phi^\dagger u_R$ and $\overline{q_L}\phi d_R$ are forbidden.

The similarity between the two mass relations $m_d m_t \approx m_c^2 \leftrightarrow m_u m_b \approx m_s^2$ suggests that the mass matrices M_{up} and M_{down} share a common structure. Such a desired feature arises naturally provided $Q_{ij}^{up} \leftrightarrow Q_{ij}^{down}$ under $\alpha \leftrightarrow -\beta$, $a \leftrightarrow -b$. The latter, to be referred to as the Up-Down symmetry, is violated of course (spontaneously) by $| \ x \ | \ll | \ y \ |$ and $| \ n \ | \ll | \ m \ |$. The various Q-charges get restricted by the Light-Heavy symmetry. The latter, being manifest via $Q_{ij} = Q_{ji}$, dictates $\alpha = b$ and $\beta = a$. In turn, the underlying $U(1)_Q$ appears to be axial, a feature long ago recognized [5] as a vital ingredient for flavor-chiral family grand unification.

To analyze the interplay of the Q-charges, let us focus attention on Q^{up}_{ij} (Q^{down}_{ij} is obtained via $a \leftrightarrow -b$). The information collected so far can be summarized by

$$Q^{up}_{ij} = \left(\begin{array}{ccc|cc|c} . & . & . & . & . & b \\ \hline . & . & . & . & b & . \\ . & . & . & a & p_1 & q_1 \\ \hline . & . & a & . & . & . \\ . & b & p_2 & . & q_2 & . \\ b & p_3 & . & . & . & . \end{array} \right) , \tag{7}$$

where $p_1 = p_2$, and p_3, q_1, q_2 have yet to be determined. Once they are specified, owing to $Q_{ij} \equiv Q_i + Q_j$, all the left over matrix elements of Q_{ij} get fixed (for example, $Q_{26} = Q_{25} - Q_{35} + Q_{36} = b - p_1 + q_1$). The locations of $q_{1,2}$ signal the block-mixing entries (see our earlier discussion) in the corresponding M^{up}_{ij}, and we have allowed for the option $p_i \neq 0$ in case the heavy mass scale M also turns out to be spontaneously generated. To determine the p's and q's as linear combinations of a and b, we first note the crucial restriction due to the fact that u_{iL} and d_{iL} form weak iso-doublets. The latter requires $Q(u_{iL}) = Q(d_{iL})$, and hence, $Q(t_L) - Q(c_L) = Q_{16} - Q_{26} = b - (b + q_1 - p_1)$ and $Q(c_L) - Q(u_L) = Q_{25} - Q_{45} = b - (a + q_2 - p_2)$ must stay invariant under $a \leftrightarrow -b$. Few other requirements such as $p_3 = \pm p_1$ (to have M entries in all p_i locations), $p_i \to \pm p_i$ under $a \leftrightarrow -b$ (to allow for the same M entries in the down sector), and $p_1 + p_2 - q_2 \neq a, q_1$ (to fully distinguish the right handed fermions), complete the picture. Altogether, we derive $p_1 = p_2 = -p_3 = -\frac{1}{2}(a + b)$ and $q_1 = q_2 = -b$.

The completion of the Q^{up}_{ij} matrix allows us, by virtue of eq.(5), to finally extract the charges of the fermions themselves. We have already arranged for $Q_{Right}(a, b) = -Q_{Left}(a, b)$ and $Q_{down}(a, b) = Q_{up}(-b, -a)$, so we only need to specify the charges of the left-handed up-quarks. We derive

$$
\begin{array}{rclcrcl}
Q(t_L) &=& \frac{1}{2}(4b-a) &,& Q(T_L) &=& -\frac{1}{2}b \quad , \\
Q(c_L) &=& \frac{3}{2}b &,& Q(C_L) &=& \frac{1}{2}(a-2b) \quad , \\
Q(u_L) &=& \frac{3}{2}a &,& Q(U_L) &=& -\frac{1}{2}a \quad ,
\end{array} \tag{8}
$$

and study their implications.

Actually, corresponding to the two degrees of freedom in eq.(8), two axial symmetries underlie our analysis:

I. Flavor-blind Z_3: The first symmetry has to do with $a = b$. This only distinguishes a standard fermion, for which $Q(u_{iL}) = \frac{3}{4}(a+b)$, from a seesaw fermion characterized by $Q(U_{iL}) = -\frac{1}{4}(a+b)$. But this is not necessarily in accord with $SU(2)_L * U(1)_Y$ which obviously requires $Q(u_{iL}) = Q(d_{iL})$. In other words, $\frac{3}{4}(a+b)$ must not change under $a \leftrightarrow -b$, a severe constraint which can only be satisfied provided the symmetry in hand is Z_3, such that $e^{i\frac{3}{2}(a+b)} = 1$. We note in passing that the associated anomaly vanishes.

II. Horizontal Z_5: The second symmetry has to do with $a+b=0$. It is easy to verify that the charges of the three seesaw fermions form an arithmetic sequence, namely $2Q(T_L) = Q(U_L) + Q(C_L)$. This may upset the Light-Heavy symmetry principle since in general $2Q(t_L) \neq Q(u_L) + Q(c_L)$. However, in analogy with the previous case, there exists a way out. One can easily verify that $e^{i\frac{5}{2}(a-b)} = 1$ is the desired constraint which makes $Q(u_{iL}) - Q(U_{iL})$ flavor-blind, thereby defining our horizontal Z_5 sub-group.

The discrete symmetry factors fit nicely in our overall philosophy. Reconstructing the $Z_3 * Z_5$ assignments of the scalars involved, we notice that $a, b = \frac{1}{2}(a+b) \pm \frac{1}{2}(a-b)$ whereas $\frac{1}{2}(a+b) = \frac{1}{2}(a+b) + 0 \cdot \frac{1}{2}(a-b)$, so that M (unlike x,y) is Z_5-invariant. This observation can be naturally translated into the spontaneous symmetry breaking chain

$$
Z_3 * Z_5 \xrightarrow{M} Z_5 \xrightarrow{x,y} 1 \ , \tag{9}
$$

thereby constituting a group theoretical origin for the Fermi mass hierarchy. This way, with M being spontaneously rather than explicitly generated, our model resembles, in some respects, the Majoron model [6] in neutrino physics. Giving up the Z_3-symmetry will allow M to be explicitly assigned, in analogy with the Gell-Mann-Yanagida model [3]. Furthermore, one cannot ignore the facts that the three singlet scalars share a common Z_3-charge and have their Z_5-charges form an arithmetic sequence. This suggests, if one is willing to introduce another scalar doublet (with a tiny VEV), a possible extension of the family structure to the Higgs system as well.

Adopting the so-called Yukawa universality [7] as a working hypothesis, and paying attention to additional entries as dictated by the $Z_3 * Z_5$ symmetry, we obtain

$$M_{up} = \left(\begin{array}{ccc|cc|c} 0 & 0 & 0 & 0 & 0 & m \\ \hline 0 & 0 & 0 & 0 & m & 0 \\ 0 & 0 & x^* & x & M & y^* \\ \hline 0 & 0 & n & 0 & 0 & 0 \\ 0 & y & M & 0 & y^* & 0 \\ y & M^* & y^* & 0 & 0 & 0 \end{array}\right) , \qquad (10)$$

accompanied by $M_{down} = M_{up}(m \leftrightarrow n^*, x \leftrightarrow y^*)$. Note that, without any lose of generality, M can be made real using three (physically unimportant) right-handed phases. Now, reordering the rows and columns of the above mass matrix, we can bring it to the canonical form $\begin{pmatrix} \alpha & \beta \\ \gamma & MI+\delta \end{pmatrix}$, where $\alpha, \beta = \mathcal{O}(m,n)$, $\gamma, \delta = \mathcal{O}(x,y)$, and the identity I are 3x3 matrices. This way it is easier to perform the perturbative expansion to obtain the effective low-energy mass matrix $m_{eff} = \alpha - \frac{1}{M}\beta\gamma + \frac{1}{M^2}(\beta\delta - \frac{1}{2}\alpha\gamma^\dagger)\gamma - \frac{1}{M^3}(\beta\delta^2 - \alpha\gamma^\dagger\delta - \frac{1}{2}\beta\gamma\gamma^\dagger)\gamma + ...$

$$m_{eff}^{(up)} \approx \begin{pmatrix} ny^2 & nxy^* & ny^{*2} \\ 0 & -mx & -my^* \\ 0 & -\frac{1}{2}mxy & m(1-\frac{1}{2} \mid y \mid^2) \end{pmatrix} , \qquad (11)$$

where, from this point on, x,y stand for $\frac{x}{M}, \frac{y}{M}$. The diagonalization procedure in both quark sectors confirms the spectrum as prescribed by eq.(4), and furthermore gives rise to the following Cabibbo-Kobayashi-Maskawa unitary matrix

$$V_{CKM} \approx \begin{pmatrix} 1 & \frac{n}{m}y^* - \frac{m^*}{n^*}x & x\left(\frac{m^*}{n^*}x - \frac{n}{m}y^*\right) \\ \frac{m}{n}x^* - \frac{n^*}{m^*}y & 1 & y^* - x \\ y\left(\frac{n^*}{m^*}y - \frac{m}{n}x^*\right) & x^* - y & 1 \end{pmatrix} . \qquad (12)$$

As advertised, reflecting the hybrid nature of the Fermi hierarchy as expressed by the $\mid \frac{m}{n} \mid$ enhancement, $\mid V_{us} \mid \approx \frac{m_c}{m_b}$ (roughly 0.23 at the conventional 1GeV mass scale) is significantly larger than the other 'nearest neighbor' $\mid V_{cb} \mid \approx \frac{m_s}{m_b}$ (roughly 0.04). All mixing angles, in particular the $\mid \frac{V_{ub}}{V_{cb}} \mid \approx \frac{m_d}{m_s}$ ratio, agree quite well with the experimental data, and so does the predicted CP-violating invariant phase. In fact, using Wolfenstein parameterization [8], one finds $\tan\phi_{CP} = \frac{Im(xy)}{Re(xy)-|x|^2}$ which can be arbitrarily large, as required, given $\mid x \mid \ll \mid y \mid$. Such an m, n - independent expression for ϕ_{CP} is quite intriguing, clearly indicating that the origin of CP-violation in the present model lies really beyond the electro/weak mass scale.

To summarize, we have attempted in this paper to account for the fermion mass hierarchy within the framework of the Universal Seesaw mechanism, where every standard fermion has a heavy seesaw partner. The imposition of a Light-Heavy symmetry leads to a Diophantine equation relating the number of families N to the sum of the distinct integers n_i characterizing the hierarchy. To our surprise, this equation has a unique non-trivial solution $N = 3$, and the hierarchy is necessarily geometric m_W, $m_W\epsilon$, $m_W\epsilon^2$. Using this fact, we have constructed a model for three quark families with a precisely defined symmetry between the up and down sectors. Hybrid quark mass relations play then a crucial role in deriving novel expressions for the CKM mixings [9] in terms of simple quark mass-ratios (to be contrasted with square-mass-ratios). To start with, in order to provide selection rules for the allowed entries in the mass matrix, we have invoked an additional $U(1)$ symmetry. However, consistency with the standard electro/weak theory allows only for its axial $Z_3 * Z_5$ sub-group, supporting the spontaneous breaking chain $Z_3 * Z_5 \xrightarrow{M} Z_5 \xrightarrow{x,y} 1$. Since only the ratios $\frac{x}{m}$ and $\frac{y}{M}$ survive in the low-energy regime, we have made no attempt to probe the M-scale itself. This is why the Yukawa universality has been invoked only at the working hypothesis level. However, the seesaw model tells us that a typical neutrino mass is of order $\frac{m^2}{M}$, so it is neutrino physics which is expected [10] to fix M. And finally, recalling that the reconciliation of string unification with low-energy may in fact require [11] exotic seesaw matter, we anticipate some of the ideas presented in this paper to find their origin in a grand unified theory.

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