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SUPERSYMMETRIC SU(11), THE INVISIBLE AXION, AND PROTON DECAY*

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ABSTRACT

We supersymmetrize the very attractive flavour unification model SU(11). As with other supersymmetric GUTs the gauge hierarchy problem is simplified, but we may also have observable ($\tau_p \approx 10^{33}$ yrs) proton decay. The required split multiplets are obtained by making the adjoint take a particular direction. Supersymmetry is broken softly at the TeV scale. There is a unique $U(1)_A$ symmetry, and hence there are no true Nambu-Goldstone bosons. The $U(1)_A$ is broken at the GUT scale and there result an invisible axion and neutrino masses.

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Supersymmetric theories [1] have attracted much attention recently [2-6], since they hold out the hope of a natural solution to the gauge hierarchy problem [7]. In conventional grand unified theories this problem arises from the existence of a large mass ratio of the order of 10^{13}, which requires a 'fine tuning' of parameters in each order of perturbation theory.

The original motivation for using supersymmetry to solve this problem was the fact that if supersymmetry is preserved to 1 TeV then the scalar partners of the left-handed electron doublet would be protected from acquiring a large mass. However in that case the (coloured) scalar partners of quarks would also be light. So if these fields are used to give masses to the fermions an unacceptably large proton decay rate would result. Hence the Higgs fields giving masses to low energy fermions cannot be the supersymmetry partners of the left-handed lepton doublets. Thus it was found necessary [4-6] to introduce additional Higgs supermultiplets. For example in SU(5) two Higgs quintuplets (H^a, $\bar{H}_a$, a = 1, 2, 3, 4, 5)* are needed. Both must be present in order to cancel anomalies arising from their fermionic partners and to give masses to both up and down quarks, since only superfields of a given chirality can occur in the superpotential.

Now the fermionic partners of the Higgs supermultiplets are in a real representation and can acquire a large mass; so the original motivation is lost. Thus in order to have a low mass Higgs doublet we again need a delicate cancellation. However we are not back to square one. The non-renormalisation of the superpotential by higher order corrections [7] implies that this adjustment of parameters has to be done only once. The coloured Higgs fields would remain heavy and the multiplet is split. However this is a somewhat bizarre situation from a group theoretic point of view.

Recently Dimopoulos and Wilczek [6] made a very important remark in this connection. They observed that such a split multiplet occurs naturally in an SU(7) model [8] spontaneously broken by an anti-symmetric Higgs field. This mechanism works only because the charge generator takes the exotic form

$$Q = \mathrm{diag}\,\left(-\frac{1}{3}\,,\,-\frac{1}{3}\,,\,-\frac{1}{3}\,,\,1\,,\,0\,,\,1\,,\,-1\right) \qquad (1)$$

However this model has the drawback of predicting exotic fermions such as a $Q = 5/3$ quark below 300 GeV. Unfortunately this mechanism cannot be generalized to higher rank groups such as SU(9) etc; for either asymptotic

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* $\bar{H}_a$ is <u>not</u> the complex conjugate of H^a.

freedom of QCD at 1 TeV would be lost or there would be no τ lepton doublet. Thus this mechanism occurs uniquely in SU(7).

In models having only ordinary quarks and leptons, the antisymmetric representation cannot be used for separating colour from the weak SU(2). The simplest way is to use an adjoint Higgs. Dimopoulos and Wilczek [6] commented on the possibility of the adjoint vanishing in the weak direction in SU(5). They found that it does not work because of the trace condition on the adjoint. In higher unitary groups the situation is somewhat different though one adjustment of parameters is still necessary. We will comment on this later after motivating the use of SU(11).

The SU(11) model first introduced by Georgi [9] has the following desirable features:

(i) It unifies flavours with non repeating fermion representations. This gives the possibility of calculating mixing angles in terms of quark masses.

(ii) The model has exactly 45 chiral fields at low energy.

(iii) Below 10^{15} GeV SU(11) behaves exactly like SU(5) apart from neutrino oscillations.

There are 561 chiral fermions in SU(11). This fact has usually been taken to be an unattractive feature of this model. However all $SU_C(3) \times SU(2) \times U(1)$ real fermions, are removed at grand unification mass scale and only the needed three generations survive. There is no intermediate mass scale between 1 TeV and 10^{15} GeV. In this sense SU(11) is the minimal model with 45 chiral fermions.

Thus it seems to us that if one wishes to use groups larger than SU(5), the most attractive candidate is SU(11). By supersymmetrizing this model we will retain these attractive features, whilst making the gauge hierarchy problem more tractable.

Suppose the superpotential contains the following term

$$W \sim H^{abc} H^d_c \bar{H}_{abd} \tag{2}$$

where the three indices in H and $\bar{H}$ are completely anti-symmetrized. Then the ordinary potential contains a term

$$V \sim | H^{abc} H^d_c |^2 + | H^d_c \bar{H}_{abd} |^2 + \text{(irrelevant)} \tag{3}$$

If the vacuum expectation value of the adjoint is

$$< H^b_a > \alpha \text{ diag } (x, x, x, 2, 2, -1, -1, -1, \ldots) \text{ ,} \tag{4}$$

where $x \neq -1$ and indices from 1 to 5 refer to the Georgi-Glashow SU(5), then we have a split multiplet and the Higgs weak doublets, for example H^{i67} ($i = 4, 5$), are massless.

Now if there are h Higgs doublet supermultiplets surviving to low energies then we have the following relations [5].

$$h = \frac{-18 + (2\pi/\alpha_{em}) [1 - 8\alpha_{em}/3\alpha_c]}{\ln M/M_W}$$

$$\sin^2\theta = \frac{3 + h/2 + (10 - h/3)\,\alpha_{em}/\alpha_c}{18 + h} \tag{5}$$

where M is the colour separation scale. For $\alpha_{em}/\alpha_c \simeq 0.0557$ and $h = 2$, one has $\sin^2\theta \simeq 0.226$ $M = 10^{17}$ GeV^5 and the proton lifetime is about 10^{40} years [10]. Then the only terrestrial test of grand unification is lost. If we wish to have a testable grand unification model i.e. one with $\tau_p \sim 10^{31}$ years, then in supersymmetric theories we need $h = 5$ giving $\sin^2\theta_w = 0.249$. In SU(5), to get four light doublets (i.e. SU(5) split multiplets) we need two adjustments of parameters. However in higher groups for example in SU(11) with just the one condition (4) six doublets survive from H^{abc} and $\bar{H}_{abc}$.

The most interesting feature of this model is that we can have [11] only one $U_A(1)$ global symmetry once the Higgs supermultiplet representations are fixed to be: H^a, H'^a, H^{abc}, H'^{abc}, H^{abcde}, H'^{abcde}, the corresponding barred representations, and the adjoint. In Georgi's SU(11) model the matter multiplets are Φ^{abcd} (330), Φ_{abc} (165), Φ_{ab} (55) and Φ_a (11). Note that except for the adjoint all other representations are the totally antisymmetric ones. In addition, to derive the symmetry breaking we introduce SU(11) singlet superfields S_1, S_1', S_3, S_3' and S_5, S_5'.

We consider the following superpotential $W + \Delta W$:

$$
\begin{aligned}
W = {} & \alpha S_1 (H^a \bar{H}_a - M_1^2) + \alpha' S_1' (H'^a \bar{H}'_a - M_1'^2) \\
& + \beta S_3 (H^{abc} \bar{H}_{abc} - M_3^2) + \beta' S_3' (H'^{abc} \bar{H}'_{abc} - M_3'^2) \\
& + \gamma S_5 (H^{abcde} \bar{H}_{abcde} - M_5^2) + \gamma' S_5' (H'^{abcde} \bar{H}'_{abcde} - M_5'^2) \\
& + \frac{\mu}{2} H_a^b H_b^a + \frac{\lambda_1}{3} H_a^b H_b^c H_c^a + \lambda_2 H^a H_a^b \bar{H}_b + \lambda_2' H'^a H_a^b \bar{H}'_b \\
& + \lambda_3 H^{abc} H_c^d \bar{H}_{abd} + \lambda_3' (H \rightarrow H') + \lambda_4 H^{abcde} H_e^f \bar{H}_{abcdf} \\
& + \lambda_4' (H \rightarrow H') + \lambda_5 H^{abc} H'^{def} H^{ghijk} \epsilon_{abcdefghijk} \\
& + \lambda_5' (H^{\cdots\cdots} \rightarrow H'^{\cdots\cdots}) \\
& + \lambda_6 \bar{H}_{abc} \bar{H}'_{def} \bar{H}_{ghijk} \epsilon \qquad + \lambda_6' (\bar{H}_{\cdots\cdots} \rightarrow \bar{H}'_{\cdots\cdots}) \\
& + \lambda_7 H^a H^{bcdef} H'^{ghijk} \epsilon \qquad + \lambda_7' (H^a \rightarrow H'^a) \\
& + \lambda_8 \bar{H}_a \bar{H}_{bcdef} \bar{H}'_{ghijk} \epsilon \qquad + \lambda_8' (\bar{H}_a \rightarrow \bar{H}'_a)
\end{aligned}
\tag{6}
$$

Mixed terms such as $H^a H_a^b \bar{H}'_b$ etc may also be present although we have not written them out explicitly.

$$
\begin{aligned}
\Delta W = {} & f_1 \Phi^{abcd} \Phi^{efgh} H^{ijk} \epsilon && + f_1' (H \rightarrow H') \\
& + f_2 \Phi^{abcd} \Phi_{abc} \bar{H}_d && + f_2' (\bar{H} \rightarrow \bar{H}') \\
& + f_3 \Phi^{abcd} \Phi_a \bar{H}_{bcd} && + f_3' (\bar{H} \rightarrow \bar{H}') \\
& + f_4 \Phi_{abc} \Phi_{de} H^{abcde} && + f_4' (H \rightarrow H'). \\
& + f_5 \Phi_{ab} \Phi_c H^{abc} && + f_6' (H \rightarrow H').
\end{aligned}
\tag{7}
$$

This superpotential has one and only one $U_A(1)$ symmetry. This symmetry is needed for strong CP invariance [12]. The fact that there is only one $U_A(1)$ symmetry is a good feature since otherwise Nambu-Goldstone boson(s) (exactly massless) would appear.

The charges under this $U_A(1)$ symmetry are as follows:

$$
\begin{aligned}
& \Phi^{abcd}; -1, \quad \Phi_{abc}; 9, \quad \Phi_{ab}; -5, \quad \Phi_a; 3, \\
& H^a; 8, \quad \bar{H}_a; -8, \quad H^{abc}; 2, \quad \bar{H}_{abc}; -2, \\
& H^{abcde}; -4, \quad \bar{H}_{abcde}; 4, \quad S_i; 0 \quad H_a^b; 0
\end{aligned}
\tag{8}
$$

The primed fields have the same charges as the corresponding unprimed ones.*

This global symmetry is the desired axial symmetry with respect to quarks and has a colour anomaly proportional to

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* If we did not have primed three index and five index Higgs fields we would have too many $U(1)_A$ symmetries with corresponding exactly massless Goldstone bosons. It is not essential to have H'^a $(\bar{H}'_a)$ unless one wishes to obtain the doublets from the one index symbols.

$$1(+3) + 9(-5) + 36(+9) + 84(-1) = 298 \neq 0$$

since ϕ^a has one quark, ϕ^{ab} has nine quarks etc.

The supersymmetry conditions

$$\frac{\partial W}{\partial S} = 0 \qquad \frac{\partial W}{\partial S'} = 0 \qquad \frac{\partial W}{\partial H^a_b} = 0 \tag{9}$$

yield the following quadratic equations for the elements C_a ($a = 1 \ldots 11$) of the adjoint Higgs in the vacuum after diagonalisation

$$\lambda_1 C_a^2 + \mu C_a - \frac{1}{11}(\lambda_1 M^2 + \lambda_2 M_1^2 + \ldots) = 0\,, \qquad 1 \leq a \leq 5 \tag{10a}$$

and

$$\lambda_1 C_a^2 + \mu C_a - \frac{1}{11}(\lambda_1 M^2 - \frac{5}{6}\lambda_2 M_1^2 + \ldots) = 0\,, \quad 6 \leq a \leq 11 \tag{10b}$$

Here $M^2 = \sum_{a=1}^{11} C_a^2$. The reason for the two equations is that $H^a(\bar{H}_a)$, H^{abc}, $(\bar{H}_{abc})$, H^{abcde} $\bar{H}_{abcde}$ and $H + H'$ acquire non-zero VeV's only in the SU(5) preserving directions at the grand unification mass scale.

The two equations (10a,b) give two roots each. However the sum of the roots for each equation must be the same. Together with the tracelessness and split multiplet (equation 4) conditions this uniquely fixes the adjoint direction

$$H^b_a = V(-5, -5, -5, 6, 6, -3, -3, -3, 4, 4, 4) \tag{11}$$

When we put $\beta = 0$ $M_3'^2 = 0$ in (6)* ($\beta' \neq 0$) this gives a total of 6 zero mass doublets, 3 each from H^{abc} and $\bar{H}_{abc}$.

Unfortunately this value of h gives with the conventional scenario $SU(11) \to SU(3) \times SU(2) \times U(1)$, $\sin^2\theta_W = 0.27$ and $\tau_p \sim 10^{29}$ yrs. Therefore

---

* Again the non-renormalisation of W allows us to do this. Note that $M_3'^2 = 0$ in order that $< H^{abc} > < H'^{abc} > = 0$. Thus no mass terms come from λ_5, λ_6 terms. $\beta' \neq 0$ in order that no light doublets come from H'^{abc}.

for this model (11) to work the symmetry breaking around 10^{15} GeV must be more complicated, so that the SU(2) coupling starts with a larger value and the colour separation scale is about 3 times larger. Then we may get even in a supersymmetric GUT an acceptable and detectable proton decay rate.

Now let us consider the question of obtaining split multiplets from the one index and five index Higgs fields. We may have two or four doublets in each case.

The adjoint takes one of the two forms below, under the conditions stated earlier in order to split the one index Higgs fields.

$$\begin{aligned} H^b_a &= V\,(2, 2, 2, 0, 0, 7, 7, -5, -5, -5, -5) \\ H^b_a &= V\,(1, 1, 1, 0, 0, 2, -1, -1, -1, -1, -1) \end{aligned} \tag{12}$$

If instead of putting $\beta = 0$ in (6) we put $\alpha = 0$ $\gamma = \gamma' = 0$ $\alpha' \neq 0$ we obtain two mass doublets. For $\alpha = \alpha' = \gamma = \gamma' = 0$ we obtain four.

Finally the five index Higgs multiplet may be split only when

$$H^b_a = V(-14, -14, -14, 20, 20, -5, -5, -5, -5, 11, 11). \tag{13}$$

In this case* for $\gamma = 0$, $\gamma' \neq 0$ we obtain $h = 2$ and for $\gamma = \gamma' = 0$ $h = 4$. This is perhaps the most attractive alternative.

For $h = 4$ we get $\tau_p = 10^{33}$ yrs and $\sin^2\theta_W \simeq 0.249$.

In case (11) (with $\beta = 0$) $M_3'^2 = 0$) the $SU(11) \times U_A(1)$ symmetry is broken down to SU(5) by the SU(5) singlets in H^{abcde}, $\bar{H} \ldots$, $\bar{H}^{\cdots}$, $\bar{H}' \ldots$, $H^a, \bar{H}_a$, H'^a, $\bar{H}'_a$. Once the adjoint has diagonalised as in (11) with the SU(6) $a = b \ldots 11$, split, the above Higgs fields clearly have enough independent components to break the symmetry $SU(11) \times U_A(1)$ completely to SU(5). Furthermore there is no possibility of the Peccei Quinn axion being ressurrected [13] – it remains invisible [14].

Similar considerations applied, mutatis mutandis, to the other two cases (12) and (13), lead to the same conclusion.

All these vacuum expectation values are of magnitude comparable to $< H^b_a >$, since they are related by the supersymmetry conditions. Thus there is only one high energy mass scale of 10^{15} GeV. In fact $SU(11) \times U_A(1)$ breaks directly to $SU(3) \times SU(2) \times U(1)$.

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* It is easily checked that λ_5 to λ_6' terms do not contribute to masses of the doublets coming from the five index symbol even if $<H^a>$, $<H^{abc}>$ etc are all non zero.

Let us now discuss the question of supersymmetry breaking. If we wish to break supersymmetry spontaneously we are confronted by the theorem of Dimopoulos and Georgi [5]; the scalar partner of either the u quark (ϕ_u) or the d quark (ϕ_d) is lighter than the corresponding quark at least at tree level. The resulting phenomenology is disastrous; for the u or the d quark may decay to its scalar partner via Goldstino emission.

There is of course the possibility of dynamical symmetry breaking as discussed by Witten [2], Dine et al. [3] and Dimopoulos and Raby [3]. However a realistic mechanism of dynamical symmetry breaking is still not known.

The remaining alternative is the soft breaking of supersymmetry as discussed by Dimopoulos and Georgi [5], for SU(5). We follow them in adding many supersymmetry breaking soft terms to the ordinary potential. Some of these terms raise the masses of unobserved particles to the TeV scale. Also a negative mass squared term for the boson fields in H^b_a fixes the direction of symmetry breaking and removes the vacuum degeneracy of the supersymmetric theory. A mixed mass squared matrix with positive and negative eigenvalues for the other Higgs fields is also added.

In the SU(5) case the soft symmetry breaking terms such as $H^\alpha H'_\alpha$ ($\alpha = 1 \ldots 5$) in the ordinary potential relate the phases of H^α and H'_α; they are equal in magnitude and have opposite signs. Hence the Peccei Quinn symmetry is removed since in SU(5) the phases should not cancel for the symmetry to be present. In SU(11) however the Higgs coming from $\bar{H}_a$, $\bar{H}_{abc}$ etc have opposite phases to those of H^a, H^{abc} etc in order to have the $U_A(1)$ symmetry (see eq. (8)). The soft terms, $H^a \bar{H}_a$, $H^a \bar{H}'_a$, $H^{abc} \bar{H}_{abc}$ etc. preserve this symmetry. Hence under soft supersymmetry breaking the Peccei-Quinn solution to the strong CP problem remains intact and the axion resulting from the spontaneous symmetry breaking at high mass is invisible as discussed before.

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