ELSEVIER

11 April 1996

Physics Letters B 372 (1996) 212–218

PHYSICS LETTERS B

A string model with $SU(4) \times O(4) \times [Sp(4)]_{\text{hidden}}$ gauge symmetry

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Received 25 September 1995; revised manuscript received 6 December 1995
Editor: R. Gatto

Abstract

In the four dimensional free fermionic formulation of the heterotic string, a semi-realistic $SU(4) \times SU(2)_L \times SU(2)_R$ model is proposed with three fermion generations in $(4,2,1) + (\bar{4},1,2)$ representations. The gauge symmetry of the model breaks to the standard gauge group using a Higgs pair in the $(4,1,2) + (\bar{4},1,2)$ representations. The massless spectrum includes exotic fractionally charged states with non-trivial transformation properties under part $(Sp(4))$ of the non-abelian 'hidden' symmetry. Finally there is a mirror pair in $(4,2,1), (\bar{4},2,1)$ allowing the possibility for an identical running of $g_{4,L,R}$ couplings between the string and $SU(4)$ breaking scales. This is of crucial importance for a successful prediction of the weak mixing angle. Potential shortcomings and problems of the construction are analysed and possible solutions are discussed.

One of the most challenging and interesting issues in strings [1] is to construct realistic models [2–10], consistent with the low energy theory. Most of the attempts in this direction [2,3,5,6] have been concentrated in constructions of string models based on level-one ($k = 1$) Kac–Moody algebras. At the $k = 1$ level, several obstacles have appeared: First, unified models based one these constructions do not contain Higgs fields in the adjoint or higher representations, therefore, traditional Grand Unified Theories (GUTs), like $SU(5)$ and $SO(10)$ could not break down to the Standard Model. Attempts to overcome this difficulty led to constructions which needed only small Higgs representations to break the symmetry [3,5]. [1] A second difficulty [11] that was encountered within the $k = 1$ Kac–Moody models was the appearance of fractionally charged states other than the ordinary Quarks, in the particle spectrum. Such states – unless they become massive at the string scale – usually create problems in the low energy effective theory. Indeed, the lightest fractionally charged particle is expected to be stable. In particular, if its mass lies in the TeV region, then the estimation of its relic abundances [12] contradicts the upper experimental bounds by several orders of magnitude. This problem can in principle be solved by constructing models containing a hidden gauge group which becomes strong at an intermediate scale and confines the fractional charges into bound states [13].

Finally, from the technical point of view, the greatest difficulty in these constructions is to obtain a three generation unified or partially unified model, which at the same time retains the successful low energy predic-

[1] more recent attempts [9,10] to overcome this difficulty have led to $SO(10) \times SO(10)$ or $SU(5) \times SU(5)$ product groups, where the $SO(10)$ or $SU(5)$ are realized directly at level 1.

PII S0370-2693(96)00059-7

tions of the supersymmetric GUT's. In fact, we know that using the Higgs and fermion content of the minimal supersymmetric standard model, the three gauge couplings $g_{1,2,3}$ of the standard gauge group attain a common value at a scale $M_{\rm GUT} \sim 10^{16}$ GeV. However in strings, the unification point ($M_{\rm string}$) is not an arbitrary parameter: it is a calculable quantity from the first principles of the theory and at the one loop level is found to be around two orders of magnitude larger that $M_{\rm GUT}$, $M_{\rm string} \sim 0.5 g_{\rm string} \times 10^{18}$ GeV. String threshold corrections [14] which can also be computed in terms of quantities related to the heavy string modes, do not bring closer these two scales. The consistency of string unification and low energy values of gauge couplings can be arranged if suitable extra matter representations and proper intermediate gauge group breaking steps are included.

A partially unified group which fulfills the basic requirements [5] is based on the Pati-Salam [15] gauge symmetry $SU(4) \times SU(2)_L \times SU(2)_R$. The symmetry can break down to the standard model gauge group without using adjoint or any higher representations. Color triplets and Higgs doublets arise in different representations, thus the model is free from doublet-triplet splitting complications, as the triplets become massive from simple trilinear couplings. There are no dangerous proton decay mediating gauge bosons, thus the $SU(4)$ breaking scale can be lower than the GUT scale predicted by other rival unified groups. Furthermore, a recent non-renormalisable operator analysis [16] of its supersymmetric version, has shown quite remarkable features on the fermion mass matrices [17,18], which provide a strong motivation to study the string derived model in more detail. The renormalisation group analysis of the string version has already been studied in detail in many papers, taking into account GUT, supersymmetric and string threshold corrections [19-21]. It was shown that it is possible to obtain the correct range of the low energy parameters while having two different scales (a string $M_{\rm string} \sim 10^{18}$ GeV and a "GUT" $SU(4)$ gauge breaking around $(10^{15} - 10^{16})$ GeV) provided there is an intermediate scale $\sim 10^{10}$ GeV where some "exotic" states acquire their masses. This was necessary to compensate for the splitting of the three standard model coupling constants, caused by the different evolution of the g_L, g_R, g_4 gauge couplings in the range $M_{\rm string} - M_{\rm GUT}$. However, a more natural way to achieve unification of the standard model gauge couplings at $\sim 10^{16}$ GeV, is to include suitable representations which enforce the same (or even approximately similar) running of the g_L, g_R, g_4 couplings between $M_{\rm string} - M_{\rm GUT}$ [16].

In the present work, we wish to present an alternative version of the string model based on a different $b_{1,2,3}$ subset of basis vectors. This new construction offers some rather interesting features with respect to its predecessor: First, the fractionally charged states appear now with non-trivial transformation properties under a hidden gauge group (namely $Sp(4)$). Although this is not probably enough to confine the fractional states at a rather high scale, the above construction can be viewed as an example how to proceed for a more realistic model. Second, due to a symmetric appearance of the $L - R$-parts of the various representations in this model, it is in principle possible to obtain almost equal values of the g_L, g_R, g_4 couplings after their running down to $M_{\rm GUT}$.

Before we proceed to the derivation of the string model, in order to make clear the above remarks we briefly start with the basic features of the supersymmetric minimal version. Left and right handed fermions (including the right handed neutrino) are accommodated in the $(4,2,1)$, $(\bar{4},1,2)$ representations respectively. Both pieces form up the complete $16^{\rm th}$ representation of $SO(10)$. The symmetry breaking down to the standard model occurs in the presence of the two standard doublet Higgses which are found in the $(1,2,2)$ representation of the original symmetry of the model. (The decomposition of the latter under the $SU(3) \times SU(2) \times U(1)$ gauge group results to the two Higgs doublets $(1,2,2) \to (1,2,\frac{1}{2}) + (1,2,-\frac{1}{2})$.) The $SU(4) \times SU(2)_R \to SU(3) \times U(1)$ symmetry breaking is realized at a scale $\sim 10^{15-16}$ GeV, with the introduction of a Higgs pair belonging to $H + \bar{H} = (4,1,2) + (\bar{4},1,2)$ representations.

The asymmetric form of the Higgs fourplets with respect to the two $SU(2)$ symmetries of the model, causes a different running for the $g_{L,R}$ gauge couplings from the string scale down to $M_{\rm GUT}$. The possible existence of a new pair of representations with $SU(2)_L$-transformation properties (as suggested in [16]) which become massive close to $M_{\rm GUT}$, could adjust their running so as to have $g_L = g_R$ at $M_{\rm GUT}$. Moreover, a relatively large number (n_D) of sextet fields ($n_D \sim 7$) remaining in the massless spectrum down to $M_{\rm GUT}$, would also result to an approximate equality of the above with g_4 coupling. Obviously, the equality of the three gauge couplings $g_{4,L,R}$ at the $SU(4)$ breaking scale $M_{\rm GUT}$, is of great importance. In practice, this means that the three standard gauge couplings $g_{1,2,3}$ start running from $M_{\rm GUT}$ down to low energies, with the same initial condition. Thus, choosing $M_{\rm GUT} \sim 10^{16}$ GeV, we are able to obtain the correct predictions for $sin^2\theta_W$ and $a_3(m_W)$. As a matter of fact, the intermediate gauge breaking step gives us one more free parameter (namely $M_{\rm GUT}$), thus having obtained the desired string spectrum we are free to choose its value in order to reconcile the high string scale $M_{\rm string}$ with the low energy data.

With the above observations in mind, we will attempt to obtain a variant of the $SU(4)\times O(4)$ model which pretty much satisfies the above requirements. The subset of the first five basis vectors we are using in our construction, including the $(1,S)$ sectors are the

$$\begin{aligned}
1 &= \{\psi^\mu,\chi^{1\ldots6},(y\bar y)^{1\ldots6},(\omega\bar\omega)^{1\ldots6};\bar\Psi^{1\ldots5}\bar\eta^{123}\bar\Phi^{1\ldots8}\}\\
S &= \{\psi^\mu,\chi^{1\ldots6},0,\ldots,0,\ 0,\ldots,0\ ;0,\ldots,0\}\\
b_1 &= \{\psi^\mu,\chi^{12},\ (y\bar y)^{3456},0,\ldots,0\ ;\bar\Psi^{1\ldots5}\bar\eta^{1}\}\\
b_2 &= \{\psi^\mu,\chi^{34},\ (y\bar y)^{12},\ (\omega\bar\omega)^{56}\ ;\bar\Psi^{1\ldots5}\bar\eta^{2}\}\\
b_3 &= \{\psi^\mu,\chi^{56},\ (y\bar y)^{1234},0,\ldots,0\ ;\bar\Psi^{1\ldots5}\bar\eta^{3}\}
\end{aligned}\tag{1}$$

All world sheet fermions appearing in the vectors of the above basis are assumed to have *periodic* boundary conditions. Those not appearing in each vector are taken with *antiperiodic* ones. We follow the standard notation used in Refs. [3,5,6]. Thus, $\psi^\mu,\chi^{1\ldots6},(y/\omega)^{1\ldots6}$ are real left, , $(\bar y/\bar\omega)^{1\ldots6}$ are real right, and $\bar\Psi^{1\ldots5}\bar\eta^{123}\bar\Phi^{1\ldots8}$ are complex right world sheet fermions. In the above, the basis element S plays the role of the supersymmetry generator as it includes exactly eight left movers. $b_{1,2}$ elements reduce the $N=4$ supersymmetries successively into $N=2,1$. Furthermore, the above set breaks the original symmetry of the right part down to an $SO(10)$ gauge group corresponding to the five $(\bar\Psi^{1\ldots5})$ complex world sheet fermions while all chiral families at this stage belong to the $16^{\rm th}$ representation of the $SO(10)$. Note here the difference of the third basis element b_3 with the one used in previous constructions [3,5,6]. To reduce further the $SO(10)$ symmetry to the desired $SO(6)\times O(4)$ gauge group, we introduce the basis elements $b_4=\{(y\bar y)^{126},(\omega\bar\omega)^{126};0,\ldots,0\}$, $b_5=\{(y\bar y)^{136},(\omega\bar\omega)^{136};\ 0,\ldots,0\}$ and the vector

$$\alpha=\{0,\ 0,\ldots,0,\ (y\bar y)^3,\ (\omega\bar\omega)^3\ ;\ \bar\Psi^{123}\bar\eta^{123}\bar\Phi^{1\ldots6}\}\tag{2}$$

These three vectors complete our basis for the model under consideration. In particular, the vector α breaks the original gauge group to the following symmetry:

$$\begin{aligned}
&[SO(6)\times SO(4)]_{\rm obs.}\times U(1)^3\\
&\quad\times[SO(12)\times Sp(4)]_{\rm Hidden}
\end{aligned}\tag{3}$$

$SO(6)\sim SU(4)$ corresponds to the three complex fermions $\bar\Psi^{123}$, while $\bar\Psi^{45}$ generate the $O(4)\sim SU(2)_L\times U(2)_R$ part of the observable gauge symmetry. $SO(12)$ corresponds to $\bar\Phi^{1\ldots6}$ while $SO(5)\sim Sp(4)$ to $\bar\omega\bar\Phi^{78}$. We have introduced subscripts to denote the *observable* and *Hidden* part of the symmetry. A well known feature of these constructions is the appearance of various $U(1)$ factors (three in the present case) which act as a family symmetry [22] between the generations. As we will see soon, the fractionally charged states in the observable sector belong also to the $4=\bar 4$ representations of the $Sp(4)\sim SO(5)$. The particular content of the model depends also on the choice of the specific set of the projection coefficients $c\left[{b_i \atop b_j}\right]=e^{i\pi c_{ij}}$. In order to guarantee the existence of $N=1$ space time supersymmetry, we choose $c\left[{S \atop b_j}\right]=1$ for $j=1,2,3$, while for the other coefficients one possible choice is $c\left[{a \atop a}\right]=c\left[{b_i \atop b_i}\right]=1$ for $i=4,5$, $c\left[{b_j \atop b_j}\right]=-1$ for $j=1,2,3$ and $c\left[{b_i \atop b_j}\right]=-1, j>i$, while all the others are fixed by the modular invariance constraints. The resulting matrix of the (exponent) coefficients c_{ij} appears in the Table 1. The upper (lower) element b_i, (b_j) of the coefficient $c\left[{b_i \atop b_j}\right]$, corresponds to the b_i-row (b_j-column) of the projection coefficient matrix exponents c_{ij} in Table 1.

We start first by presenting the spectrum with the representations which are going to be interpreted as fermion generations and $SU(4)$ breaking Higgses. Fermion generations arise from $b_{1,2,3}$ sectors appearing in symmetric representations under the $SO(6)\times O(4)$ symmetry. Thus it makes no difference which of the two resulting representations of $b_{1,2,3}$ will accommodate the left or right components of the fermion generations. The choice of the assignment however, is crucial for the Higgs fourplets which are

Table 1
The exponent factors c_{ij} of the projection coefficient matrix.

	1	S	b_1	b_2	b_3	b_4	b_5	α
1	1	0	1	1	1	1	1	0
S	0	0	0	0	0	1	1	1
b_1	1	1	1	1	1	1	1	1
b_2	1	1	1	1	1	1	1	1
b_3	1	1	1	1	1	1	1	1
b_4	1	1	1	1	1	0	1	1
b_5	1	1	1	1	1	1	0	1
α	1	1	1	1	1	1	1	0

not symmetric under the two $SU(2)$'s. Thus, starting with one of the two possible choices the sectors which provide with the fermion generations and possible $SU(4)$ breaking Higgses are[2]

$$\begin{array}{lll}
b_1: & F_1 & = (4,2,1)_{(-1/2,0,0)}; \\
 & \bar{F}_1 & = (\bar{4},1,2)_{(1/2,0,0)} \\
b_2: & F_2 & = (4,2,1)_{(0,-1/2,0)}; \\
 & \bar{F}_2 & = (\bar{4},1,2)_{(0,1/2,0)} \\
b_3: & F_3 & = (4,2,1)_{(0,0,-1/2)}; \\
 & \bar{F}_3 & = (\bar{4},1,2)_{(0,0,-1/2)} \\
b_2+b_4: & F_{24} & = (4,2,1)_{(0,1/2,0)}; \\
 & \bar{F}_{24} & = (\bar{4},1,2)_{(0,-1/2,0)} \\
b_3+b_4+b_5: & \bar{F}_{345} & = (\bar{4},2,1)_{(0,0,1/2)}; \\
 & \bar{F}'_{345} & = (\bar{4},1,2)_{(0,0,-1/2)}
\end{array} \tag{4}$$

The above representations of the observable sector transform trivially under the hidden gauge group. However, they all appear charged under the three $U(1)$ factors corresponding to $\bar{\eta}_1, \bar{\eta}_2, \bar{\eta}_3$ world-sheet fermions. These charges are denoted with the three indices in the above representations. $F_{1,2,3}, \bar{F}_{1,2,3}$ can accommodate the three generations, while from the $(b_2 + b_4)$ and $(b_3 + b_4 + b_5)$ sectors we get a pair of family-antifamily ($F_{24} - \bar{F}_{345}$) left-fourplets. Unfortunately, in this case the two remaining representations $\bar{F}'_{345}, \bar{F}_{24}$ cannot play the role of the two $SU(4)$ Higgses, as they are both of the type $\bar{H}_{1,2} = (\bar{4},1,2)$. More over, this spectrum apparently creates an anomaly with respect to the $SU(4)$ gauge group, since there is an excess of fourplet over anti-fourplet fields; however, there is a pair of exotic states $(4,1,1)^{n}_{(1/2,0,0)} + (4,1,1)^{\bar{n}}_{(1/2,0,0)}$ with fractional charges arising from the sector $(1+b_2+b_3+b_4+\alpha)$ which guarantee the anomaly cancellation. The novel feature of these representations here, is their non-trivial transformation under part of the hidden non-abelian gauge group. In fact they belong to the $n = \bar{n} = 4$ representation (denoted as superscript) of the $Sp(4)$ symmetry. As we will see soon, this is also true for the rest of the exotics in this construction. Provided the hidden group confines at some later stage, this allows for the possibility of forming various types of condensates. By choosing proper flat directions, such states may become massive and disappear from the light spectrum, while some of them can have the right Higgs properties so that they can be used to break the $SU(4)$ symmetry. Indeed, in order to examine this case further, in the following let us continue with the relevant representations. From the sectors $(1+b_1+b_2+\alpha)$, $(1+b_1+b_2+b_4+b_5+\alpha)$ and $(1+b_2+b_3+b_4+\alpha)$ we obtain six pairs of "exotic" doublet states $(1,1,2)^{(n/\bar{n})} + (1,2,1)^{(\bar{n}/n)}$, possessing half-integer ($\pm 1/2$) electric charges. Interestingly enough, these exotic states can in principle condense with the $(4,1,1)^{n}_{(1/2,0,0)} + (4,1,1)^{\bar{n}}_{(1/2,0,0)}$ states into the missing Higgs fourplets $H_{1,2} = (4,1,2)$ at a later scale. (Their $U(1)$-charges depend on the specific $(1,1,2)$ representations.) Thus in this way there can exist now two Higgs pairs (namely $H_{1,2} + \bar{H}_{1,2}$) where either of them can break the $SU(4)$-symmetry to the standard model. However, of crucial importance is the confinement scale M_C of the $Sp(4)$ symmetry, as it simultaneously defines the $SU(4)$ breaking scale of the observable symmetry. This can be calculated from the formula

$$M_C = M_{\rm string} \exp\{\frac{2\pi}{b_{SO_5}}(\frac{1}{\alpha_{\rm string}} - \frac{1}{\alpha_c})\} \tag{5}$$

where $b_{SO_5} = -3C_2(SO_5) + 2n_4 + n_2$ is the beta function of $SO(5)$, while $C_2(SO_5) = 3$. For two fourplet Higgses we need $n_4 = n_2 = 2$ thus $b_{SO_5} = -3$ as in the case of the $SU(3)$, which means that the confining scale is rather low. However, there are some important differences which should be mentioned. First, the initial scale where the renormalisation starts is $M_{\rm string}$ which is two orders higher than the supersymmetric unification scale $M_{\rm GUT}$. Furthermore the unified coupling $a_{\rm string}$ turns out to be larger than the common gauge coupling $a_{\rm GUT}$ in the minimal supersymmetric unification. For example in [23] it is found

[2] The second case arises by interchanging $4 \leftrightarrow \bar{4}$, $2_L \leftrightarrow 2_R$ in the above sectors, and will be commented below.

$a_{\rm string} \sim 1/20$, while $a_{\rm GUT} \sim 1/25$. Thus, in contrast to the $SU(3)$, for the $Sp(4)$ confining scale one finds $M_{Sp_4} \sim 10^7$ GeV. This scale is still rather low compared to the usual grand unification. However, in the case of the $SU(4)$ 'unification' this is not a disaster; as we have already pointed out, there are no gauge bosons mediating proton decay, thus a low energy breaking scale is not necessarily in contradiction with the low energy phenomenology. Nevertheless, it would be desirable to obtain a rather higher confinement scale close to the 'conventional' minimal supersymmetry unification point $\sim 10^{15-16}$ GeV. This of course would need a confining group with rank higher that the $Sp(4)$.

From the Neveu-Schwarz sector we get the following fields: Two Higgs fields of the type $(1,2,2)_{(0,0,0)}$ under the observable $SU(4) \times SU(2)_L \times SU(2)_R$ gauge group, and no charges under the three family-type $U(1)$ symmetries. Six sextet fields $(6,1,1)_{(\pm1,0,0)+{\rm perm.}}$ Various singlet fields $\chi^{\alpha}_{(\pm1,0,\pm1)}, \chi^{\beta}_{(\pm1,\pm1,0)}, \chi^{\gamma}_{(0,\pm1,\pm1)}$ with integer (±1) surplus $U(1)$ charges are also available. Representations with the same transformation properties but different charges under the three $U(1)$-family symmetries are obtained from the sectors $S+b_2+b_3$, $S+b_1+b_3$ and $S+b_1+b_3+b_4$. In particular, they give singlet fields analogous to those of the NS-sector but with half-integer extra $U(1)$ charges, $\xi^{\alpha}_{(\pm1/2,0,\pm1/2)}, \xi^{\beta}_{(\pm1/2,\pm1/2,0)}, \xi^{\gamma}_{(0,\pm1/2,\pm1/2)}$, and $\Sigma_{(\pm1,\pm1/2,\pm1/2)}$. In addition in the massless spectrum there exist vector representations of the hidden part of the symmetry which do not have transformation properties under the observable gauge group. Thus, each of the above three sectors gives the **12** of $SO(12)$ and **5** of $SO(5)$. The resulting three $12^{\rm th}$ irreps do not play any role in the observable world, however if the **5**'s remain massless, they can lower dangerously the confining scale. Finally, from the same sectors one gets sextet fields $D_{1,2,3} = (6,1,1)_{(0,1/2,1/2)}, (6,1,1)_{(\pm1/2,0,1/2)}$ and Higgses $h_{1,2,3} = (1,2,2)_{(0,-1/2,1/2)}, (1,2,2)_{(\pm1/2,0,1/2)}$. At least one of the latter is expected to acquire a vacuum expectation value (vev) along its two neutral components in order to give masses to fermion generations through Yukawa couplings allowed by gauge and string symmetries. Although only few couplings are expected to be present at the trilinear superpotential, there is a large variety of singlet fields possessing various $U(1)$ charges which are going to form non-renormalisable mass terms.

Let us briefly now discuss the fermion masses. Light fermions acquire their masses with the usual Higgs mechanism, when some of the $(1,2,2) \to (1,2,\frac{1}{2}) + (1,2,-\frac{1}{2})$ Higgs representations develop vevs. If we assume that below $M_{\rm GUT}$ the model behaves approximately as the minimal supersymmetric standard model, only one pair of the available electroweak Higgs doublets (or only a linear combination of them) should remain light. Then, in the trilinear superpotential, a coupling of the form $\lambda^0_{ijk} F_i \bar{F}_j h_k$ will provide with masses the fermions of the third generation, with the GUT-predictions $m^0_t = m^0_{\nu_D}$, $m^0_b = m^0_\tau$, where $m^0_{\nu_D}$ is the Dirac neutrino mass. A remarkable feature of these string models is the generic prediction that the Yukawa coupling λ^0_t responsible for the top-quark mass is large and of the same order with the common gauge coupling at the string scale, $\lambda^0_t = \sqrt{2} g_{\rm string}$, leading to a top mass of the $\mathcal{O}(180)$ GeV [23]. This is compatible with previously proposed SUSY-GUT models which predicted radiative symmetry breaking and a large top mass with a single third generation Yukawa coupling [24]. The bad $(m^0_t, m^0_{\nu_D})$ relation is handled with the "see-saw"-type relation through a term of the form $H\bar{F}_i\Phi_n \to \langle H\rangle \nu_{R_i}\Phi_n$ as described in previous works [5,18]. The rest of the entries of the fermion mass matrices are expected to fill up when non-renormalisable contributions to the superpotential are taken into account. Additional colored triplets $d^c_H, \bar{d}^c{}_H$ remaining from the $H+\bar{H}$ representations form massive states with $D_3, \bar{D}_3$ states arising from the decomposition of the sextet fields $D \to D_3 + \bar{D}_3$, through terms of the form $HHD, \bar{H}\bar{H}D$ [5]. Note that some of them could be harmless even if they get mass at a relatively low scale $\sim 10^7$ GeV provided they do not couple with the ordinary matter at the tree level.

Finally, the family $F_{24} = (4,2,1)$-antifamily $\bar{F}_{345} = (\bar{4},2,1)$ pair can become massive either at the tree level or from a higher order non-renormalisable coupling of the form $\mathcal{W} \supset \langle\Phi_i\rangle(4,2,1)(\bar{4},2,1)$, with $\langle\Phi_i\rangle \sim M_{\rm GUT}$. In fact the singlet vevs are not completely arbitrary in these constructions. From the three family type $U(1)$'s of the present model, one can define two linear combinations (say $U(1)_1 - U(1)_2 - U(1)_3, U(1)_2 - U(1)_3$) which are anomaly free, while the remaining orthogonal combination remains anomalous. The latter is broken by the Dine-Seiberg-Witten mechanism [25] in which a potentially large supersymmetry-breaking D-term is generated, by the vacuum expectation value of the dilaton field. To avoid this situation, one has to choose a D- and F-flat direction in the scalar potential by assigning proper vevs to some of the scalar fields. The natural scale of these singlet vevs turns out to be $M_{\rm string} \geq \langle \Phi_i \rangle \geq M_{\rm GUT}$.

Let us finally analyse the alternative accommodation of the fermion generations and higges under the observable symmetry. This can be easily obtained by interchanging $4 \leftrightarrow \bar{4}$ and $2_L \leftrightarrow 2_R$ in the relevant sectors. The three sectors $b_{1,2,3}$ provide again the three generations. From $b_2 + b_4$ one gets $F_{24} = (4,2,1), \bar{F}_{24} = (\bar{4},1,2)$ while $b_3 + b_4 + b_5$ gives $F_{345} = (4,1,2), F'_{345} = (4,2,1)$. (Of course the $U(1)$ charges are not affected.) Thus, now the Higgs fourplets $H + \bar{H}$ needed to break the $SU(4)$ symmetry are contained in the $(b_2 + b_4)$ and $(b_3 + b_4 + b_5)$ sectors. In fact we can now identify $H \equiv F_{345} = (4,1,2)$ and $\bar{H} \equiv \bar{F}_{24} = (\bar{4},1,2)$. It is possible however that a detailed phenomenological analysis of the model would require some linear combinations of F_i's and $\bar{F}_i$'s to be interpreted as the $SU(4)$ breaking Higgses of the model. Thus, in this case the Higgs particles are not formed by condensates, therefore the 'GUT' scale is not related to the confinement scale. We may choose then $M_{\rm GUT} \sim 10^{15-16}$ GeV and obtain a renormalisation group running of the gauge couplings as described above. The present accommodation however, creates a new problem; the two remaining pieces of $(b_2 + b_4)$ and $(b_3 + b_4 + b_5)$ sectors, have the same transformation properties with the left handed fermion generations. These two remaining $(4,2,1)$ states are rather difficult to become massive. However, it is possible that after the $SU(4)$ breaking the resulting colored triplets and doublets may combine with their conjugate partners arising from the composite states $(\bar{4},2,1)$ (which now transform as anti-fourplets under the interchange $4 \leftrightarrow \bar{4}$) through non-renormalisable terms resulting in an effective mass term much lower than the scale $M_{\rm GUT}$.

The above model is not of course a fully realistic model for the low energy theory. However, it is a rather interesting improvement of a previous version which was based on the same gauge symmetry. Its advantages with respect to the old version can be briefly summarized in the following points: Fractionally charged states transform non trivially under a hidden gauge group (namely $Sp(4)$) which forces them to form bound states. Specific composite states can play the role of the Higgses which break the $SU(4) \times SU(2)_R$ symmetry while the most of the remaining hopefully may combine in various terms with other fields into relatively heavy massive states escaping detection by the present experiments. The main drawback of this construction is that the $Sp(4)$ group falls rather short to confine these charges at a suitably high scale. A novel feature of this construction of the model is also the choice of the vectors $b_{1,2,3}$ which are different from those already used in the flipped $SU(5)$ [3] and standard model [6] constructions. Since the previous $SU(4)$ model has been pretty much similar to the flipped $SU(5)$ we think that the three basis vectors $b_{1,2,3}$ used here can also offer new possibilities for these constructions which are worth exploring.

The author would like to thank the CERN (Geneva) and Ecole Polytechnique (Palaiseau) HEP-Theory Groups for kind hospitality. This work was partially supported by CEC project CHRX-CT93-0132.

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