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Supercompositeness from superstrings

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Abstract. String Unified Models based on the k = 1 level of the Kac-Moody Algebra, predict the existence of "exotic" new states which carry fractional electric charges. We analyse the possibility of considering these "exotics" as preonic matter which can be used to form the families and the gauge group breaking higgs scalars. It is proposed that such a formation may occur provided that these states transform non-trivially under a non-Abelian gauge group with a relatively large rank in order to confine them at a sufficiently large scale. Such a situation is natural in string derived unified models, since the role of the confining group can be played by (part of) the Hidden symmetry. As an example, we present a string derived toy model based on the $SU(4) \times SU(2)_L \times SU(2)_R$ Pati-Salam gauge group.

One of the most attractive and interesting issues in strings, is the construction [1] of realistic models which are consistent with the low energy phenomenology. Most of the attempts in string model building [2–7] have been concentrated in constructions of string models based on level-one (k=1) Kac-Moody algebras. In the search for the realistic string derived models, two main obstacles have appeared:
(i) Unified models based on these constructions do not contain 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 the construction of models where the gauge group needs only small Higgs representations to break [2, 4].
(ii) The appearance of fractionally charged states, other than the ordinary Quarks, is unavoidable [8] in the k=1 Kac-Moody constructions. Such states, unless they become massive at the String scale, they 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 [9] contradicts the upper experimental bounds by several orders of magnitude. It has been proposed that this problem can in principle be solved by constructing models containing a hidden gauge group which becomes strong at an intermediate scale to confine the fractional charges into bound states [2, 10].

In this work, we would like to explore an alternative scenario: Since the fractionally charged states are generic [8] in k=1 level, it might be possible that in particular string models they could play the role of some *preonic* matter superfields which transform non-trivially under some *Hidden* gauge group. This *Hidden* group could very well play the role of the confining gauge group of the preonic fields into composite states which could be the representations containing the ordinary Quarks and Leptons.

Models with composite Quarks and Leptons have already been introduced by many people [11–15] the last two decades, either in the context of the Standard gauge group of Electroweak interactions[1] or within Grand Unified schemes. Both scenarios are well motivated in the context of Superstring k=1 constructions. Indeed, if we insist on the economy of the models derived from the String, we would feel unhappy with a large variety of representations left in the light spectrum of the effective field theory, even if we manage with a judicious choice of the parameters of the theory (moduli, flatness conditions etc) to make them massive at some intermediate scale M_C. Instead, it would appear more natural to derive an effective field theory with a relatively small number of representations.

In what follows, motivated by the appearance of such exotic states in string constructions, we will concentrate on a particular gauge group which leads to a viable low energy scenario. In particular we will explore the $SU(4) \times SU(2)_L \times SU(2)_R$ Pati-Salam symmetry assuming the existence of representations carrying fractions of the known electric charges possessed by the ordinary Quarks and Leptons. In this work, we will not discuss long standing problems which arise when trying to implement the idea of compositeness. Relevant discussions about the problems appearing in the various approaches to the compositeness may be found in the literature [11, 12, 17].

In the free fermionic four dimensional constructions [1], in principle it is possible to choose boundary conditions on

[1] Compositeness may also be combined with Technicolor and Extended Technicolor Theories [16] to produce interactions which may create dynamically the fermion masses

the world sheet fermions of the basis vectors of a particular model and project out all the integrally charged representations [2]. The fermion families, will appear then at an intermediate scale as composite fields of the preonic representations. Therefore, it would be natural to ask, if any phenomenologically viable preonic model could arise from the Plank scale physics.

The necessary fields of the minimal supersymmetric version together with their transformation properties under the PS-gauge symmetry are shown in the table

$$F_L = (4,2,1);\ \bar{F}_R = (\bar{4},1,2);\ H = (4,1,2);\quad \bar{H} = (\bar{4},1,2);$$
$$h = (1,2,2);\quad D = (6,1,1);\quad \phi_{m,0} = (1,1,1),\ m = 1,2,3$$

The superpotential of the model can be written as follows

$$\begin{aligned}\mathscr{W} &= \lambda_1 F_L \bar{F}_R h + \lambda_2 \bar{F}_R H \phi_i + \lambda_3 \phi_0 h h + \lambda_4 \phi^3 \\ &\quad + \lambda_5 H D D + \lambda_6 \bar{H}\bar{H} D + \lambda_7 \bar{F}_R \bar{F}_R D + \lambda_8 F_L F_L D \end{aligned} \tag{1}$$

The superpotential (1) includes trilinear terms with states arising only from the decomposition of the ordinary irreducible representations (irreps) of the $SO(10)$ theory. At k=1 Kac-Moody level in particular, all irreps appearing in the theory are smaller than the adjoint. For the model under consideration for example, the possible representations under $SU(4)\times SU(2)_L\times SU(2)_R$ arise from the decompositions of $\underline{16}$, $\overline{\underline{16}}$ and $\underline{10}$ of SO(10). In string constructions, however, the case is more complicated. In fact, in this particular model the effective theory gauge symmetry is based on a product of non-Abelian groups rather than on a single unified one. In the fermionic constructions for example, the model is constructed from a set of vectors whose components are phases picked up by the world-sheet fermions when parallel transported around non-contractable loops. The massless states of the theory are those surviving the projections of the various vectors onto the others. As a result, in addition to the above states, new representations may arise which are singlets under all but one of the non-Abelian factors of the symmetry of the model. Thus in addition to $(4,2,1),(\bar{4},1,2)$ one may get the "exotics" $(4,1,1),(\bar{4},1,1)$, while together with $(1,2,2)$ one also obtains $(1,2,1)$ and $(1,1,2)$. Of course such representations are not present in the ordinary $SO(10)$ irrep decompositions.

There are two ways of handling these states:

i) One can redefine the charge operator [8, 4, 18]. Indeed, in the usual string constructions the resulting "observable" gauge symmetry is accompanied by "hidden" gauge groups and a rather large number of U(1) factors. Most of the fields discussed above carry non-zero charges under the surplus U(1)'s. One then could extend the charge operator by including one or more of these U(1)'s. Such cases have been discussed in the literature but they usually lead to the wrong predictions for the weak mixing angle.

ii) As a second possibility we consider the case discussed here where the string model predicts only the "exotic" states discussed above, with non-trivial transformation properties under part of the hidden gauge group. String toy models with such properties can be easily constructed [19].

With the form of the minimal theory in mind, let us now attempt to derive it considering only *preonic* fields, assuming that the ordinary superfields are not present in the original theory. As has already been discussed, we assume that the symmetry of the observable sector is based on the gauge group $SU(4)\times SU(2)_L\times SU(2)_R$, while the fields belong also to some $N(\bar{N})$-dimensional representation of a *Hidden* gauge group. The fractionally charged states which appear in the string spectrum of these models, in the most general case, are of the following types

$$\begin{array}{llll} K_j & = (4,1,1)_{jN} & K_n^c & = (\bar{4},1,1)_{nN} \\ \bar{K}_n & = (4,1,1)_{n\bar{N}} & \bar{K}_j^c & = (4,1,1)_{j\bar{N}} \\ \alpha_{Li} & = (1,2,1)_{iN} & \alpha_{Rm} & = (1,1,2)_{mN} \\ \bar{\alpha}_{Lm} & = (1,2,1)_{m\bar{N}} & \bar{\alpha}_{Ri} & = (1,1,2)_{i\bar{N}} \end{array} \tag{2}$$

Let us explain our notations in the above fields. The numbers in the parentheses, as usually, refer to the transformation properties of the various *preonic* fields under the observable gauge symmetry of the model. The indices i,j,m,n refer to the number of the corresponding representations and run from 1 to $\mathscr{I},\mathscr{J},\mathscr{M},\mathscr{N}$ respectively. Care has been taken, so as $(4/\bar{4})$ as well as $(N/\bar{N})$ representations appear in pairs to ensure that the theory is anomaly free. The index $N(\bar{N})$ in each of the above representations refers to its transformation property under the *Hidden* gauge group [3].

We should note here, that in realistic string constructions, the fields might also carry extra $U(1)$-charges while the *Hidden* gauge group is not always simple. However, in order to make the subsequent analysis simple and model independent, we consider only a simple $SU(N)$-*Hidden* gauge group. The existence of the $U(1)$ factors would have the obvious implication of reducing the possible gauge invariant trilinear and higher order Yukawa terms of the superpotential.

Now, if we define the charge operator in the usual sense

$$\mathscr{Q} = \frac{1}{6}T_{15} + \frac{1}{2}T_L + \frac{1}{2}T_R \tag{3}$$

where $T_{15} = diagonal(1,1,1-3)$, and $T_{L,R}$=$diagonal(1,-1)$, it is clear that all the above fields carry charges which are fractions of those of ordinary Quarks and Leptons. Under the *Hidden* gauge group, they form composite states at some intermediate scale $M_U < M_C < M_{Pl}$, which may be identified with the ordinary superfields of Quarks and Leptons. The possible composite states created from the above *preonic* fields are listed in Table 1.

We observe that all the fields of the superpotential in (1) are present in Table 1. The indices $\{i,j,m,n\}$ indicate the multiplicity of each representation. Thus, in the general case considered above, one obtains $\mathscr{J}\mathscr{M}+\mathscr{I}\mathscr{N}$ left handed fields F_L and an equal number of right handed representations $\bar{F}_R$. These representations are going to accommodate the known fermion families of quarks and leptons and their superpartners. Note however, that the above are accompanied by $\mathscr{I}\mathscr{J}+\mathscr{M}\mathscr{N}$ "anti-left" $\bar{F}_L$, and "anti-right" F_R fields. The rest of the composite spectrum includes $\mathscr{I}^2+\mathscr{M}^2$ higgses in (1,2,2), $2\mathscr{J}\mathscr{N}$ (6,1,1)-sextets and an equal number of (10,1,1) irreps. The new feature is the appearance

[2] In the $SU(4)\times SU(2)_L\times SU(2)_R$ model this is rather obvious. Indeed the charge operator is a combination of diagonal generators of all the group factors. Therefore, in the fermionic language for example, we may extend the basis by adding "phase-vectors" until all representations transform non-trivially under only one of the non-Abelian groups. Many of these representations do not belong to any known "GUT" multiplets and possess charges which are fractions of those of the ordinary fermions

[3] Sextets under $SU(4)$ can also arise at this level

Table 1.

$(F_L)_{jm}$	$=$	$K_j\bar{\alpha}_{Lm}$	$= (4,2,1)$	;	$(F_R)_{ji}$	$=$	$K_j\bar{\alpha}_{Ri}$	$= (4,1,2)$
$(\bar{F}_L)_{ni}$	$=$	$K_n\alpha_{Li}$	$= (\bar{4},2,1)$	;	$(\bar{F}_R)_{nm}$	$=$	$K_n\alpha_{Rm}$	$= (\bar{4},1,2)$
$(\bar{F}_L)_{nm}$	$=$	$K^c_n\bar{\alpha}_{Lm}$	$= (\bar{4},2,1)$	;	$(\bar{F}_R)_{ni}$	$=$	$K^c_n\bar{\alpha}_{Ri}$	$= (\bar{4},1,2)$
$(F_L)_{ij}$	$=$	$K^c_j\alpha_{Li}$	$= (4,2,1)$	;	$(F_R)_{jm}$	$=$	$K^c_j\alpha_{Rm}$	$= (4,1,2)$
D_{jn}	$=$	$K_j\bar{K}_n$	$= (6,1,1)$	;	D_{nj}	$=$	$K^c_n\bar{K}^c_j$	$= (6,1,1)$
T_{jn}	$=$	$K_j\bar{K}_n$	$= (10,1,1)$	;	T_{nj}	$=$	$K^c_n\bar{K}^c_j$	$= (10,1,1)$
Σ_{jj}	$=$	$K_j\bar{K}^c_j$	$= (15,1,1)$	;	$\bar{\Sigma}_{nn}$	$=$	$K^c_n\bar{K}_n$	$= (15,1,1)$
Φ_{jj}	$=$	$K_j\bar{K}^c_j$	$= (1,1,1)$	;	$\bar{\Phi}_{nn}$	$=$	$K^c_n\bar{K}_n$	$= (1,1,1)$
h_{ii}	$=$	$\alpha_{Li}\bar{\alpha}_{Ri}$	$= (1,2,2)$	;	h_{mm}	$=$	$\alpha_{Rm}\bar{\alpha}_{Lm}$	$= (1,2,2)$
Δ_{Lim}	$=$	$\alpha_{Li}\bar{\alpha}_{Lm}$	$= (1,3,1)$	;	Δ_{Rmi}	$=$	$\alpha_{Rm}\bar{\alpha}_{Ri}$	$= (1,1,3)$
Φ'_{im}	$=$	$\alpha_{Li}\bar{\alpha}_{Lm}$	$= (1,1,1)$	;	$\bar{\Phi}'_{mi}$	$=$	$\alpha_{Rm}\bar{\alpha}_{Ri}$	$= (1,1,1)$

of $\mathscr{J}^2+\mathscr{N}^2$ adjoint representations (15,1,1) of SU(4) and $\mathscr{T}\mathscr{M}$ (1,3,1) and (1,1,3) as well as $2\mathscr{T}\mathscr{M}+\mathscr{J}^2+\mathscr{N}^2$ neutral singlets.

In order to avoid the appearance of "antifamilies" in the light spectrum, they should combine with equal number of families and receive mass at a high scale. If we demand r generations to remain in the light spectrum, then we should have $\#F_L-\#\bar{F}_L=r$, and an equal number of right partners. This requirement leads to the equation $(\mathscr{M}-\mathscr{T})\times(\mathscr{J}-\mathscr{N})=r$, which is satisfied for various choices of $\mathscr{T},\mathscr{J},\mathscr{M}$, and $\mathscr{N}$.

Thus, let us distinguish some simplified cases:

• $\mathscr{M}=0$. In this case, the above requirements for $r=3$ lead to the condition $(\mathscr{N}-\mathscr{J})\times\mathscr{T}=3$. An acceptable case for three generations would be $\mathscr{J}=1$, $\mathscr{N}=2$ and $\mathscr{T}=3$. In this case one obtains 6 F_L's, 3 $\bar{F}_L$'s, 3 F_R's and 6 $\bar{F}_R$'s. In addition, there are 4 D's, 9 h's and 5 singlets. In order to remain with the minimal spectrum of the superpotential of (1), three pairs $\bar{F}_L+F_L$ should become massive through some effective superpotential term of the form $<\Phi>\bar{F}_LF_L$. As far as the right representations are concerned, one pair $\bar{F}_R+F_R$ should be interpreted as the higgses $\bar{H}+H$ which break the $SU(4)$ symmetry, while the remaining two additional pairs should become massive in the same way as the left fields. In a similar way, we can give superheavy masses to any other additional representations like sextets and doublets.

• $\mathscr{J}=0$. In this case the condition reads $(\mathscr{T}-\mathscr{M})\times\mathscr{N}=3$. We may further choose $\mathscr{T}=4$, $\mathscr{M}=1$, $\mathscr{N}=1$ or $\mathscr{T}=2$, $\mathscr{M}=1$, $\mathscr{N}=3$. For the second case for example we obtain the following spectrum: $6(F_L+\bar{F}_R)$, $3(\bar{F}_L+F_R)$, $5(1,2,2)$ LR-higgses 9 $(15,1,1)$-higgs in the adjoint of SU(4) and 2 $(1,3,1)+(1,1,3)$ pairs accompanied by 18 neutral singlet fields. There are no sextet fields but usually they arise naturally in the original spectrum of the particular string model.

The remarkable feature in this spectrum is the presence of two types of higgses which both can break the $SU(4)$ symmetry. This fact gives various possibilities of symmetry breaking patterns which will be described in a future publication.

In the following we present a string toy example [19] based on the Pati-Salam (PS) observable gauge symmetry and a specific hidden interaction which can play the role of confining gauge group, previously referred as $SU(N)$. This will enable us to give a more detailed description of the suggested approach of the previous sections. We will work in the free fermionic formulation of the four dimensional superstring using free world-sheet fermions which pick up phases [1] when parallel transported around the string. A specific choice of phases – consistent with the modular invariance constraints – for all world-sheet fermions of the left (supersymmetric) and right (non-sypersymmetric) sectors of the string, define a 'phase' vector. The specific choice of the boundary conditions, will determine the exact gauge group of the theory and the specific transformation properties of the preonic fields we are using. More precisely, these vectors break simultaneously the original high string ($SO(44)$) symmetry to the observable PS-gauge group and the hidden part. In particular, the supermultiplets of the observable gauge group are formed from the fractionally charged states which transform as 4 or $\bar{4}$ representations of an $SU(4)$ hidden gauge group. Although this is not a fully realistic model, it serves as an example where the basic features described so far in the field theory approach of the previous section, as well as the basic ingredients of the model, can be found in the string spectrum of $k=1$ models. The rank of the $SU(4)_{Hidden}$ group is smaller than the one needed to confine the preonic matter at the conventional GUT scale of $10^{16}GeV$. In the most optimistic case, the $SU(4)$ confining scale is of the order $10^{12}GeV$, i.e. $3-4$ orders smaller than the conventional GUT scale. In models with PS-gauge symmetry such a low scale is not disastrous. In fact, there are no gauge bosons mediating proton decay and the SU(4) breaking scale can be low enough provided this is consistent with the low energy values of gauge couplings. This is indeed the case revealed in several recent renormalization group analyses [20, 21] of models based on the PS- gauge group. Note however that a realistic model should at least have $SU(5)$ or a higher rank symmetry as a confining group.

Let us now present the basis of 'phase vectors' generating the preonic matter discussed above, at the string level. We denote the 18 real free fermions of the supersymmetric left-moving sector with $\chi^{1\ldots6},y^{1\ldots6},\omega^{1\ldots6}$. For the right moving 22 complex fermions we use the notation $\bar{\Psi}^{1\ldots5},\bar{\Phi}^{1\ldots9}$ and $\bar{z}^{1\ldots8}$. Now a particular basis is generated by assuming consistent [1] boundary conditions on the world-sheet fermions of the two sectors of the heterotic string. In our case, we assume six vectors of boundary conditions which form a group under addition modulo 2. The basis is the following:

$$\begin{array}{lll}
1 &=& \{\psi^\mu,\ \chi^{1\ldots6},\ y^{1\ldots6},\ \omega^{1\ldots6};\ \bar{\Psi}^{1\ldots5}\bar{\Phi}^{1\ldots9}\bar{z}^{1\ldots8}\}\\
S &=& \{\psi^\mu,\ \chi^{1\ldots6},\ 0,\ldots,0;\ 0,\ldots,0\ \ 0,\ldots,0\}\\
b_1 &=& \{\psi^\mu,\ \chi^{12},\ y^{3456},\ 0,\ldots,0;\ \bar{\Psi}^{123}\bar{\Phi}^{123}\underbrace{\bar{z}^{1\ldots8}}_{1/2}\}\\
b_2 &=& \{\psi^\mu,\ \chi^{34},\ y^{1256},\ 0,\ldots,0;\ \bar{\Psi}^{45}\bar{\Phi}^{123}\underbrace{\bar{\Phi}^{78}}_{1/2}\underbrace{\bar{z}^{1\ldots6}}_{-1/2}\bar{z}^{7}\}\\
b_3 &=& \{00,\ 0\ldots0,\ y^{1\ldots6},\ \omega^{1\ldots6};\ \bar{\Phi}^{78}\}\\
\zeta &=& \{00,\ 0\ldots0,\ 0,\ldots,0,\ 0,\ldots,0;\ \bar{\Phi}^{123}\bar{\Phi}^{678}\bar{z}^{12}\}
\end{array} \quad (4)$$

All world sheet-fermions appearing in the vectors – except those underlined with $\pm1/2$ – possess periodic boundary conditions, while those not appearing in a particular vector are antiperiodic. The symmetry breaks down to the following product group

$$[SU(4)\times SU(2)_L\times SU(2)_R]_{Observable}$$

$$\times[\{SU(4)^2\}_C \times SU(4)' \times SU(2)' \times SU(2)'' \times U(1)^6]_{Hidden} \quad (5)$$

where the *Observable* and the *Hidden* sectors of the symmetry are denoted with subscripts. The particular content of the model depends on the choice of the specific set of the projection coefficients $c\begin{bmatrix} b_i \\ b_j \end{bmatrix}$. One possible choice is $c\begin{bmatrix} 1 \\ 1 \end{bmatrix} = c\begin{bmatrix} \zeta \\ \zeta \end{bmatrix} = -1$, $c\begin{bmatrix} b_1 \\ b_1 \end{bmatrix} = c\begin{bmatrix} b_2 \\ b_2 \end{bmatrix} = -1$, $c\begin{bmatrix} b_i \\ \zeta \end{bmatrix} = -1$, $c\begin{bmatrix} S \\ b_2 \end{bmatrix} = -1$, while $c\begin{bmatrix} i \\ j \end{bmatrix}$=+1 for the remaining with $i \leq j$ in the order of appearance in (4). All the others are fixed from the modular invariance constraints.

The spectrum which arises is listed below, where the quantum numbers refer to the *Observable* sector groups, the two $\{SU(4)^2\}_C$ groups and the 6 $U(1)$s.

$$\begin{array}{lll} b_1: & K_1^c & = (\bar{4},1,1)(4,1)_{(0,0,1/2,1,1/4,1/4)} \\ 3b_1: & K_2 & = (4,1,1)(4,1)_{(0,0,-1/2,-1,-1/4,-1/4)} \\ b_2: & \alpha_{1L} & = (1,2,1)(4,1)_{(1/2,0,-1/2,-1,1/2,0)} \\ 3b_2: & \bar{\alpha}_{2L} & = (1,2,1)(4,1)_{(-1/2,0,1/2,1,-1/2,0)} \\ 2b_1+b_2: & \bar{\alpha}_{1R} & = (1,1,2)(4,1)_{(1/2,0,1/2,1,0,-1/2)} \\ 2b_1+3b_2: & \alpha_{2R} & = (1,1,2)(4,1)_{(-1/2,0,-1/2,-1,0,1/2)} \end{array} \quad (6)$$

An equal number of fractionally charged states transforming as 4 or $\bar{4}$ under the second of the $\{SU(4)^2\}_C$, differing only in the U(1) factors, arises if we add the vector ζ to the above sectors. The above fields are accompanied by singlet fields ξ_i, ζ_i which arise from the Neveu-Schwarz (N-S) and ζ sectors and two $(6,1,1)+(1,2,2)$ representations from the $2b_1+3b_2+\zeta$ sector. Finally there are four $SU(4)' \times SU(2)'$ representations transforming as $Q=(4,2)'$, $\bar{Q}=(\bar{4},2)'$.

All the resulting representations of the observable sector have double multiplicity while they transform as the 4 or $\bar{4}$ under one of the two $SU(4)_C$ groups. In particular, fermion like condensates arise from the combinations $F_{1L}=\bar{K}_2\alpha_{1L}$ and $F_{1R}=K_1^c\bar{\alpha}_{1R}$. Similarly one can obtain the rest of the spectrum presented in Table 1. Unfortunately, the number of preonic states in this toy example does not meet the conditions put previously in order to obtain the right number of generations and higgs fields, at least not directly at the string level. However, in general, this fact does not exclude the possibility of finding a flat direction where some of the singlet fields get non-zero vevs and make massive the superfluous preonic fields through trilinear couplings of the superpotential shown in the Appendix. Thus, our string toy example, although not a realistic one, it is deductive with respect to what one should expect from a string derived spectrum. For example, families (and other fields) are distinguished from each other by different $U(1)$ factors accompanying the PS symmetry. Thus the family constructed above has quantum numbers $F_{1L}=(4,2,1)_{(\frac{1}{2},0,-1,-2,\frac{1}{4},-\frac{1}{4})}$ and $F_{1R}=(\bar{4},1,2)_{(\frac{1}{2},0,1,2,\frac{1}{4},-\frac{1}{4})}$. The appearance of U(1) symmetries is an encouraging fact as it may generate the desired fermion mass hierarchy.

Let us finally discuss how we reach the $N=1$ supersymmetry. The element S of the above basis, with exactly 8 left movers, plays the role of supersymmetry generator in the fermionic construction. The subset $\{1,S,1+\zeta\}$ defines an $N=4$ space-time supersymmetric model, while the introduction of the rest of the basis vectors break successively the $N=4$ supersymmetries to $N=2$ and $N=1$. In the class of models we are dealing with, the confinement scale M_C is assumed to be no less than the GUT scale, i.e., the $SU(4)$-breaking scale. Thus, at the scale M_C we are left with an $N=1$ supergravity model based on PS-observable gauge symmetry. Now, the only consistent way to break $N=1$ local supersymmetry is spontaneously via the super Higgs mechanism. No matter what the mechanism of supersymmetry breaking is (dynamical breaking [22], gaugino condensation [23], coordinate dependent compactifications [24] etc) a natural hierarchy $m_{3/2} \ll M_{Planck}$ should be generated, where $m_{3/2}$ is the gravitino mass which sets the scale of supersymmetry breaking. This in turn implies that there should exist a class of string effective supergravities where certain conditions are satisfied like absence of fine tuning, vanishing cosmological constant up to $\mathcal{O}(m_{3/2}^4)$ corrections etc. It has been shown [25] that such conditions can be met particularly in the four dimensional string fermionic constructions we are dealing with in this work. In particular, supersymmetry breaking via gaugino condensation can be shown [25] to exist in examples for superpotentials with non-trivial dependence on the dilaton field S, with a well behaved positive-semi-definite potential. However, as the the Hidden gauge group of this kind of models is rich, a dynamical supersymmetry breaking scenario [22] is quite possible here [26]. In such models, the messengers could be states of a Hidden gauge group. In the present case, the Q, $\bar{Q}$ representations of the $SU(4)' \times SU(2)'$ Hidden symmetry do not couple to the ordinary matter representations, whilst they form trilinear couplings $\Phi_i Q\bar{Q}$ with the gauge singlet fields $\Phi_i=\xi_i,\zeta_i$ of the N-S and ζ sectors, in the tree level superpotential. Due to the presence of the massless matter multiplets in the spectrum, the Hidden gauge group $SU(4)'$ may be strongly interacting at a relatively low scale Λ. Indeed, the beta function $b_4'=-12+2n_4=-4$ in this case, while the scale is given by

$$\Lambda = M_{string} exp\{\frac{2\pi}{b_4'}(\frac{1}{\alpha_{string}}-\frac{1}{\alpha_\Lambda})\} \quad (7)$$

Thus, for $M_{string} \sim 5\times10^{17}GeV$ and $\alpha_{string}=1/24$, α_Λ becomes ~ 0.2 at $\Lambda \sim 5\times10^4 GeV$. Now, one of the singlet fields may obtain a non-vanishing vev along the scalar and F- components. As a result, gaugino masses arise at one loop, being directly proportional to the scale Λ

$$m_i \sim \frac{\alpha_i}{4\pi}\Lambda \quad (8)$$

This shows that the scale Λ is of the desired order for a dynamical symmetry breaking scenario.

In conclusion, in this note we have examined the possibility of obtaining low energy effective gauge models using only representations which arise at k=1 level of string constructions, with gauge symmetry based on the $SU(4)\times SU(2)_L\times SU(2)_R$ Pati-Salam (PS) gauge group. We have argued that even though the adjoint or higher representations are absent from the spectrum of such string derived models, there are still various possibilities of obtaining a consistent low energy phenomenology. In particular, in constructions based on the aforementioned PS-gauge symmetry and k=1

Kac-Moody level, we have seen that a viable low energy phenomenological theory can be derived in the following ways:

i) One can use the standard PS-representations $(4,2,1)$, $(\bar{4},1,2)$ to accommodate the three fermion families, while the standard gauge symmetry is obtained after the spontaneous breaking of the PS symmetry with a minimum number of Higgses sitting in $(4,1,2)$ + $(\bar{4},1,2)$ representations [4]. The standard model can break with the use of the two standard doublets found in the (1,2,2) of the PS-symmetry. These standard representations are accompanied by the "exotics" $(4,1,1)$, $(\bar{4},1,1)$ and $(1,2,1)$, $(1,1,2)$ which carry fractional electric charges, and should form massive states at a rather high scale to avoid phenomenological problems.

ii) As a second possibility, we have argued in this paper that the above "exotics" may arise with non-trivial transformation properties under a "hidden" gauge group with sufficient rank, in order to confine to integral charged states at a high scale. It has been shown that the resulting condensates can have the correct transformation properties to accommodate quark, lepton and higgs fields and reproduce the model of case i). In addition, new symmetry breaking patterns can be obtained as it is possible now to accommodate higgs fields in the adjoint representation. The models of case ii) are reminiscent of supersymmetric composite models, proposed sometime ago [13, 12, 14].

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