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N=1 Superstrings with Spontaneously Broken Symmetries¹

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ABSTRACT

We construct N=1 chiral superstrings with spontaneously broken gauge symmetry in four space time dimensions. These new string solutions are obtained by a generalized coordinate dependent Z_2 orbifold compactification of some non-chiral five dimensional N=1 and N=2 superstrings. The scale of symmetry breaking is arbitrary (at least classically) and it can be chosen hierarchically smaller than the string scale $(\alpha')^{-1/2}$.

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String solutions with spontaneously broken gauge symmetries were recently constructed^[1] by means of a generalized coordinate dependent compactification, the analog on strings of the Scherk-Schwarz mechanism^[2] used in field theory.

In a recent work we presented the partition functions of string solutions with $N=2$, $N=4$, $N=6$, $N=8$ supersymmetries^[3] and broken gauge symmetries, and we proved their multiloop modular invariance and factorization properties (cluster decomposition) on higher genus surfaces^[4]. The mass spectrum of these models depends upon several continuous parameters, which may eventually be determined through quantum corrections.

The existence of string solutions, based on a multi-loop modular invariant partition function, which depend upon continuous parameters is not surprising. In fact, it has been realized that the same solutions correspond to an underlying two dimensional superconformal theory with world sheet gauge bosons coupled to some world sheet complex fermions,^[4] or alternatively, to certain types of Thirring strings. The world sheet gauge boson and Thirring interpretation suggest that the same models may be obtained using free world sheet fermions which satisfy generalized boundary conditions with arbitrary phases.^[5]

Whatever the interpretation is, the knowledge of the one-loop partition functions allows to show that these continuous parameters are the order parameters for gauge symmetry breaking in the target space.

The previously proposed breaking mechanism^[1] can only be used to break symmetries in models with an even number of space-time supersymmetries.

This is a major obstruction for the application of this mechanism to realistic models with chiral families which allow the presence of $N=1$ supersymmetry at most. However, this can be avoided by a suitable twist of the extra (fifth) dimension. In the present note we show that by using twisted boundary conditions for the fifth coordinate, (Z_2 -orbifold compactification)^[6,7] we introduce a suitable projection which reduces the number of supersymmetries down to one. The resulting string model contains chiral fermion and gauge symmetry is spontaneously broken at an arbitrary scale. Schematically we have:

$$\left(\begin{array}{c} N=1, D=5 \\ \text{superstring} \end{array} \right) \xrightarrow{\tau(z_5)/z_5} \left(\begin{array}{c} N=1, D=4 \\ \text{spontaneously} \\ \text{broken superstring} \end{array} \right) \quad (1)$$

where by $T(z_5)/z_5$ we denote the z_5 -dependent orbifold compactification. Notice that the simple z_5 -torus compactification gives in four dimensions a $N=2$, $D=4$ superstring while the z_5 -dependent torus compactification $T(z_5)$ gives a spontaneously broken version of the same theory.^[1] Notice also that in the standard orbifold compactification,^[6] T/z_5 , the resulting $N=1$, $D=4$ superstring has an exact space time $N=1$ supersymmetry.

In general one may start from any $N=1$, $D=5$ superstring; a variety of those models can be easily constructed by using the technique of ref. [8,9,10] where all internal string degrees of freedom were described by free world sheet fermions.^[10] We may classify them in two groups, according to the number of world sheet supersymmetries (i) heterotic models with $(1,0)$ two dimensional supersymmetry, and (ii) type II models with $(1,1)$ supersymmetry. In both cases to realize (non-linearly) the world sheet supersymmetry,^[11,12] we use free fermions which are in the adjoint representation of $H_L = [SO(3)]^5$ or in the adjoint of $H_R = SU(4)$. The choice of H_L in both cases follows from the requirement of five dimensional target space time supersymmetry, when the boundary conditions of the world sheet fermions are chosen periodic and antiperiodic.^[11,9] This restriction is avoided when more general boundary conditions are chosen.^[10] For definiteness we restrict our attention to simple models based on $(1,0)$ world sheet supersymmetry. Our mechanism, however, is general and can be applied to more complicated models.

In the heterotic construction that we examine here, the left moving degrees of freedom (analytic two dimensional fields) are given, in the light cone gauge, by:

$$\begin{aligned} \partial_+ X^\mu(z), \psi^\mu(z), X^I(z), y^I(z), w^I(z) \\ \mu = 3, 4, 5, \quad I = 2, 3, 4, 5, 6 \end{aligned} \quad (2)$$

X^μ are the bosonic string coordinates, ψ^μ are their left moving superpartners and (X^I, y^I, w^I) are left moving fermions triplets under $SO(3) \subseteq H_L, I = 2, 3, 4, 5, 6$.

The right moving degrees of freedom (anti-analytic fields) are:

$$\begin{aligned} \partial_- X^\mu(\bar{z}), \quad \phi^A(\bar{z}) \\ \mu = 3, 4, 5, \quad A = 1, 2, 3, \dots 42. \end{aligned} \quad (3)$$

X^μ are the bosonic string coordinates and ϕ^A are right moving fermions.

We will choose the $N=1$ $D=5$ superstring defined by the basis elements,^[9]

$$F = \{\psi^\mu, \psi^5, \chi^2, \dots, \chi^6, y^2, \dots, y^6, w^2, \dots, w^6, \phi^1, \dots, \phi^{42}\},$$

$$S = \{\psi^\mu, \psi^5, \chi^2, \chi^3, \dots, \chi^6\},$$

$$b_1 = \{\psi^\mu, \psi^5, \chi^2, y^3, y^4, y^5, y^6, \phi^1, \dots, \phi^{16}\}. \quad (4)$$

As was shown in ref. [8,9], the F, S, b_1 define two dimensional fermion number projections, $\frac{1}{2}(1 \pm (-)^F)$, $\frac{1}{2}(1 \pm (-)^S)$ and $\frac{1}{2}(1 \pm (-)^{b_1})$. The Hilbert space of the theory contains states of the type $R^a N^b F^c$ where a is a set of fermions with periodic boundary conditions (Ramond) and the rest (F^c) is given antiperiodic boundary conditions (Neveu-Schwarz). The physical states are those with definite fermion number $(-)^F$, $(-)^S$ and $(-)^{b_1}$. The choice of those fermion numbers is fixed by the requirement of multiloop modular invariance and factorization.^[9]

The (F, S) model in five dimensions defines the $N=2$ $SO(42)$ superstring.^[9] The introduction of b_1 , (F, S, b_1) defines a $N=1$ $D=5$ model with gauge group $SO(16) \times SO(26)$. The simple torus reduction give rise to a $N=2$ $D=4$ model while with the standard orbifold compactification one may obtain $N=1$, $D=4$ superstrings based on the gauge groups:

$$SO(10) \times SO(6) \times SO(5 + 8k) \times SO(21 - 8k), \quad k = 0, 1, 2 \quad (5)$$

The action of Z_2 on the world sheet fields which defines the above orbifold models is

$$\begin{aligned} \hat{g}\varphi &= \varphi\hat{g} \quad \text{for } \varphi \notin b_2, \\ \hat{g}\varphi &= -\varphi\hat{g} \quad \text{for } \varphi \in b_2, \quad \hat{g}^2 = 1 \end{aligned} \quad (6)$$

where b_2 is the following set of fields:

$$b_2 = \left\{ X_5; \psi^\mu, \chi^3, \chi^5, y^2, y^4, y^6, \underbrace{\phi^1, \dots, \phi^{10}}_{10}, \underbrace{\phi^{17}, \dots, \phi^{21+8k}}_{5+8k} \right\}, \quad k = 0, 1, 2. \quad (7)$$

The choice of b_2 (the Z_2 action) is consistent with the global existence of the world sheet supercurrent T_F

$$\hat{g}T_F = -T_F\hat{g}, \quad (8)$$

where

$$T_F = \partial_\alpha X^\mu \psi_\mu + \partial_\alpha X^5 \psi_5 + \sum_{I=2}^6 \chi^I y^I w^I \quad (9)$$

The right moving fermion content in b_2 has been chosen in order to generate a $SO(10)$ gauge group which admits families in the complex representation 16. As we will see later on, the choice of the χ^I and y^I fields in b_2 is also restricted when the T/Z_2 orbifold is generalized to $T(x_3)/Z_2$ by a consistency requirement.

For the of simplicity we start first with the partition function of the (F, S, b_1) model compactified on a torus T

$$Z(\tau) = Z_B \sum_{\substack{\vec{a}_{L,R} \\ \vec{b}_{L,R}}} Z_{m,n}^s(\tau).$$

$$C \left[\begin{matrix} \vec{a}_L & \vec{a}_R \\ \vec{b}_L & \vec{b}_R \end{matrix} \right] Z_L \left[\begin{matrix} \vec{a}_L \\ \vec{b}_L \end{matrix} \right] (0|\tau) Z_R \left[\begin{matrix} \vec{a}_R \\ \vec{b}_R \end{matrix} \right] (0, \tau) \quad (10)$$

where Z_B is the contribution of the X^3, X^4 bosonic string coordinates

$$Z_B = \tau_I^{-1} |\eta(\tau)|^{-4} \quad (11)$$

$Z_{m,n}$ is the winding number dependent torus partition function

$$Z_{m,n} = R(\tau_I)^{-1/2} |\eta(\tau)|^{-2} \exp - \frac{\pi R^2}{\tau_I} |m - \tau n|^2 \quad (12)$$

(in our normalization the radius of the X_5 torus is $\sqrt{2}R$)

$Z_L \left[\begin{matrix} \vec{a}_L \\ \vec{b}_L \end{matrix} \right]$ and $Z_R \left[\begin{matrix} \vec{a}_R \\ \vec{b}_R \end{matrix} \right]$ are the contribution of the left and right moving fermions given as product of θ -functions^[13] with characteristics depending on the spin structure assignment. Their explicit form is^[14]

$$\begin{aligned} Z_L \left[\begin{matrix} \vec{a}_L \\ \vec{b}_L \end{matrix} \right] (U_L, \vec{V}_L|\tau) &= e^{-i\pi U_L} \prod_{i=1}^{N_L} \frac{\theta \left[\begin{matrix} a_L^i \\ b_L^i \end{matrix} \right] (V_L^i|\tau)}{\eta(\tau)} \\ \text{and} \\ Z_R \left[\begin{matrix} \vec{a}_R \\ \vec{b}_R \end{matrix} \right] (U_R, \vec{V}_R|\tau) &= e^{+i\pi U_R} \prod_{i=1}^{N_R} \frac{\theta \left[\begin{matrix} a_R^i \\ b_R^i \end{matrix} \right] (V_R^i|\tau)}{\eta(\tau)} \end{aligned} \quad (13)$$

where $2N_L(2N_R)$ is the number of left (right) fermionic degrees of freedom, and variables $U, \tilde{V}_L, \tilde{V}_R$ have been introduced for future use.

The coefficients $C \begin{bmatrix} \tilde{a}_L & \tilde{a}_R \\ \tilde{b}_L & \tilde{b}_R \end{bmatrix}$, satisfy some relations ensuring the modular invariance of eq.(10)^[6,9] as well as the correct factorization of their multi-loop extension.^[9]

As was shown in ref.[1] the partition function of all coordinate dependent torus compactifications is

$$Z(e_L^i, e_R^i | \tau) = Z_B \sum_{\substack{\tilde{a}_{L,R} \\ \tilde{b}_{L,R}}} \sum_{m,n} Z_{m,n}^B(\tau) \cdot C \begin{bmatrix} \tilde{a}_L & \tilde{a}_R \\ \tilde{b}_L & \tilde{b}_R \end{bmatrix} \cdot Z_L \begin{bmatrix} \tilde{a}_L \\ \tilde{b}_L \end{bmatrix} (U_L, \tilde{V}_L | \tau) Z_R \begin{bmatrix} \tilde{a}_R \\ \tilde{b}_R \end{bmatrix} (U_R, \tilde{V}_R | \bar{\tau}) \quad (14)$$

where $U_{L,R}, \tilde{V}_{L,R}$ are winding number depending charges defined as

$$\begin{aligned} \tilde{V}_L &= \tilde{e}_L(m - \tau n), U_L = \tilde{e}_L \tilde{e}_L n(m - \tau n) \\ \tilde{V}_R &= \tilde{e}_R(m - \bar{\tau} n), U_R = \tilde{e}_R \tilde{e}_R n(m - \bar{\tau} n). \end{aligned} \quad (15)$$

The parameters $e_{L,R}^i$ are arbitrary and define the strength of the gauge symmetry breakdown. We would like now to extend T and Tx_5 to orbifold compactifications^[6,7] T/z_2 and $T(x_5)/z_2$ in order to obtain chiral models. The Z_2 action is already defined in eqs.(6,7) through the set of fields b_i . As we will see immediately, the choice of fermions in b_2 (or, in general, any other set of fermions f) must satisfy constraints dictated by the presence of non-zero charges e_L^i, e_R^i . In fact, the Z_2 orbifold^[6,7] is obtained from the torus, by identifying the points x_5 and $-x_5$, while the fermions in b_2 define a fermion number operator $(-)^f$. The Z_2 action projects out some states according to their fermion number $(-)^f$. As was shown in ref.[1] by using the trace representation of $\theta \begin{bmatrix} a \\ b \end{bmatrix} (V | \tau)$ functions^[13] the $T(x_5)$ partition function of eq.(14) can be expressed as a sum of products of the form

$$\text{Tr} \mathcal{P}_{F,q} L_0 \tilde{q}^{L_0}, \quad q = e^{2\pi i \tau},$$

where $\mathcal{P}_{F_i}$ are fermion number projections, and $L_0, \tilde{L}_0$ the e_L^i and e_R^i dependent Virasoro operators given by:^[1]

$$\begin{aligned} L_0 + \tilde{L}_0 &= \frac{1}{2R^2} (m + \tilde{e}_L \tilde{Q}_L - \tilde{e}_R \tilde{Q}_R - 1/2(\tilde{e}_L^2 - \tilde{e}_R^2)n)^2 \\ &+ \frac{1}{2} R^2 n^2 - (\tilde{e}_L \tilde{Q}_L + \tilde{e}_R \tilde{Q}_R)n + \frac{1}{2} (\tilde{e}_L^2 + \tilde{e}_R^2)n^2 + L'_0 + \tilde{L}'_0 \\ L_0 - \tilde{L}_0 &= mn + L'_0 - \tilde{L}'_0 \end{aligned} \quad (16)$$

Here, $L'_0, \tilde{L}'_0$ are all the terms not depending on the winding number (m, n) . Q_L^i and Q_R^i are left moving and right moving $U(1)$ charge operators.^[1]

From this equation and noticing that the state with winding number (m, n) is sent into that with winding number $(-m, -n)$ by $x_5 \rightarrow -x_5$, and since states conjugated under the orbifold Z_2 projection must have the same mass, we see that the fermion number operator $(-)^f, f \subseteq b_2$, must anticommute with all charges Q_L^i and Q_R^i . In other words, $(-)^f$ should act like a charge conjugation on all the world sheet fermions having e_L^i or e_R^i different from zero. So, we get the additional consistency condition

$$[Q_{L,R}^i, (-)^f]_+ = 0 \quad \forall i \quad \text{with } e_{L,R}^i \neq 0 \quad (17)$$

Naturally, $Q_{L,R}$ must also commute with the other fermion number projections F, S, b_1 .

A possible right moving charge is, for instance, that which rotates ϕ^{10} and ϕ^{11} .

Left moving charges must also commute with the world-sheet supercurrent, beyond fulfilling eq.(17).

We may now define the $T(x_5)/z_2$ orbifold model.

As in the T/z_2 case,^[6,7] the $T(x_5)/z_2$ partition function is written as sum of 4 contributions:

$$Z_0(+, +), Z_0(+, -), Z_0(-, +) \text{ and } Z_0(-, -).$$

Their explicit form is:

$$(a) \quad Z_0(+, +) = Z(\tilde{e}_L, \tilde{e}_R | \tau) \text{ as in eq.(14)}.$$

$$(b) \quad Z_0(+, -) = Z_B Z^s(+, -) \sum_{\substack{\vec{a}_{L,R} \\ \vec{b}_{L,R}}} C \left[\begin{array}{c} \vec{a}_L \\ \vec{b}_L + \vec{g}_L \end{array} \middle| \begin{array}{c} \vec{a}_R \\ \vec{b}_R + \vec{g}_R \end{array} \right] \quad (22)$$

$$Z_L \left[\begin{array}{c} \vec{a}_L \\ \vec{b}_L + \vec{g}_L \end{array} \middle| \begin{array}{c} \vec{a}_R \\ \vec{b}_R + \vec{g}_R \end{array} \right] (0|\vec{\tau})$$

$$(c) \quad Z_0(-, +) = Z_B Z^s(-, +) \sum_{\substack{\vec{a}_{L,R} \\ \vec{b}_{L,R}}} C \left[\begin{array}{c} \vec{a}_L + \vec{h}_L \\ \vec{b}_L \end{array} \middle| \begin{array}{c} \vec{a}_R + \vec{h}_R \\ \vec{b}_R \end{array} \right] \quad (23)$$

$$Z_L \left[\begin{array}{c} \vec{a}_L + \vec{h}_L \\ \vec{b}_L \end{array} \middle| \begin{array}{c} \vec{a}_R + \vec{h}_R \\ \vec{b}_R \end{array} \right] (0|\vec{\tau})$$

$$(d) \quad Z_0(-, -) = Z_B Z^s(-, -) \sum_{\substack{\vec{a}_{L,R} \\ \vec{b}_{L,R}}} C \left[\begin{array}{c} \vec{a}_L + \vec{h}_L \\ \vec{b}_L + \vec{g}_L \end{array} \middle| \begin{array}{c} \vec{a}_R + \vec{h}_R \\ \vec{b}_R + \vec{g}_R \end{array} \right] \quad (24)$$

$$Z_L \left[\begin{array}{c} \vec{a}_L + \vec{h}_L \\ \vec{b}_L + \vec{g}_L \end{array} \middle| \begin{array}{c} \vec{a}_R + \vec{h}_R \\ \vec{b}_R + \vec{g}_R \end{array} \right] (0|\vec{\tau}) \quad (18)$$

where:

$$\vec{g}_{L,R} = (g_{L,R}^1, g_{L,R}^2, \dots), \quad \vec{h}_{L,R} = (h_{L,R}^1, h_{L,R}^2, \dots)$$

and

$$h_{L,R}^i, g_{L,R}^i = 1 \text{ if } i \in b_i; \quad h_{L,R}^i, g_{L,R}^i = 0 \text{ if } i \notin b_i. \quad (19)$$

In eqs.(20), $Z^s(+, -)$, $Z^s(-, +)$ and $Z^s(-, -)$ are the contributions of the twisted X^s bosonic coordinate.^[7]

$$Z^s(+, -) = \left| q^{\frac{1}{2}} \prod_n (1 + q^n) \right|^{-2}$$

$$Z^s(-, +) = \left| q^{-\frac{1}{2}} \prod_n (1 - q^{n+\frac{1}{2}}) \right|^{-2} \quad (20)$$

$$Z^s(-, -) = \left| q^{\frac{1}{2}} \prod_n (1 + q^{n+\frac{1}{2}}) \right|^{-2}$$

In terms of these quantities the 1-loop partition function reads:

$$Z_0^{g=1}(\vec{e}_L, \vec{e}_R|\tau) = \frac{1}{2} [Z_0(+, +) + Z_0(+, -) + Z_0(-, +) + Z_0(-, -)] \quad (21)$$

At this point a remark is in order: at any loop the partition function is given as

$$Z_0^g = 2^{-g} [Z_{torus}^g + Z_{twist}^g] \quad (22)$$

where Z_{torus}^g and Z_{twist}^g factorize separately on Riemann surfaces "pinched" along separating nodes. This property follows from the factorization of the Z_2 orbifold partition function, examined in ref.[7] and well known properties of the fermionic chiral determinants.^[14]

In order to determine the particle content of the theory we must perform a Poisson resummation on the m winding number in eq.(21). Only the $Z_0(+, +)$ part has a winding number dependence and is identical to that of ref.[1], (see eq.(18a)), so we can immediately obtain the L_0 and $\bar{L}_0$ operators in the various sectors, as well as the associated Z_2 projections. In the untwisted sector L_0 and $\bar{L}_0$ are exactly as in the $T(x_5)$ torus given in eq.(16).

$$(L_0 + \bar{L}_0)_{UTS} = (L_0 + \bar{L}_0)_{T(x_5)} \quad (23)$$

The twisted sector is even simpler since it does not depend on the value of the charges $e_{L,R}^i$. The spectrum remains identical to that of $e_{L,R}^i$ i.e. of the (T/z_2) orbifold. A comment is now in order. In the spectrum of the $T(x_5)$ compactification, before the Z_2 projection, for every state $|q, R\rangle$ with charge $Q_i = q_i$ and all the other quantum numbers, denoted here by R , equal, there was a state $|-q, R\rangle$ with the same mass. This degeneracy is removed by the introduction of $(-)^{b_2}$ projection. $(-)^{b_2}$ acts differently on states with charges $q \neq 0$ or $q = 0$. When $q \neq 0$, $(-)^{b_2}$ selects a particular linear combination of $|q, R\rangle$ and $|-q, R\rangle$ ($|-q, R\rangle + |q, R\rangle$) or $(|-q, R\rangle - |q, R\rangle)$ according to whether R is even or odd under $(-)^{b_2}$. When $q = 0$, the $(-)^{b_2}$ projection eliminates all the states with $(-)^{b_2}|0, R\rangle = -|0, R\rangle$. The fact that states of non-zero charge (and also non-zero winding numbers) appear in pairs of opposite $(-)^{b_2}$ parity implies

$$Z_0(+, -) = \text{Tr} \left((-)^{b_2} q^{\frac{L_{UTS}}{2}} \right) \left| \begin{array}{c} Q_i=0 \\ m, n=0 \end{array} \right|$$

$$= \text{Tr} \left((-)^{b_2} q^{\frac{L_{UTS}}{2}} \right) \left| \begin{array}{c} Q_i \neq 0 \\ m, n \neq 0 \end{array} \right| \quad (24)$$

Eq.(24) shows that $Z_0(+, -)$ is an index, counting only the states without charges or winding numbers. This fact explains why $Z_0(+, -)$ is independent

of the value of the charges $e_{L,R}^i$ and guarantees the correct particle interpretation of the generalized $T(x_s)/Z_2$ orbifold model. Indeed, due to that property,

$$\frac{1}{2} (Z_0(++)) + Z_0(+-) = Tr \frac{1}{2} (1 + (-)^{b_s}) q^{L_{0s}^{UTS}} \bar{q}^{L_{0s}^{UTS}} \quad (25)$$

and

$$\frac{1}{2} (Z_0(-+)) + Z_0(--)) = Tr \frac{1}{2} (1 + (-)^{b_s}) q^{L_{0s}^{LTS}} \bar{q}^{L_{0s}^{LTS}} \quad (26)$$

The above equations show the consistency of the $(-)^{b_s}$ projection. Comparing the orbifold model T/Z_2 with the generalized $T(x_s)/Z_2$, we see that the states which are physical in the limit $e_L^i = e_R^i = 0$ remain physical for every value of e_L^i, e_R^i without any additional states being introduced. This proves that $T(x_s)/Z_2$ is the spontaneously broken version of the T/Z_2 model. The mass shifts induced by $e_L^i, e_R^i \neq 0$ is given by eq.(1b) where now Q is replaced by $\frac{1}{2}(1+h)Q$, and $h = +1$ in the untwisted sector while $h = 1$ for the twisted one.

In order to understand the symmetry breaking pattern of the $T(x_s)/Z_2$ theory it is sufficient to examine the masses of the light states (massless in the limit $e_{L,R}^i \neq 0$). For these states one finds (eq.(16)) with $Q \rightarrow \frac{1}{2}(1+h)Q$)

$$M^2|_{m=n=0} = \frac{1}{4R^2} (1+h) (\vec{e}_L \vec{Q}_L - \vec{e}_R \vec{Q}_R)^2 \quad (26)$$

To exemplify eq.(26), we keep all the e_R and e_L equal to zero, except the right charge e which parametrizes the rotation of ϕ^{10} and ϕ^{11} . By this choice, the gauge groups $G_{k=0} = SO(10) \times SO(6) \times SO(5) \times SO(5) \times SO(21)$ and $G_{k=1} = SO(10) \times SO(6) \times SO(13) \times SO(13) \times SO(13)'$ (remember that $k=2$ is equivalent to $k=0$) are broken to:

$$\begin{aligned} G_{k=0} &= SO(9) \times SO(5) \times SO(5) \times SO(21) \times SO(2) \\ G_{k=1} &= SO(9) \times SO(5) \times SO(13) \times SO(13)' \times SO(2) \end{aligned} \quad (27)$$

We will examine first the twisted sector of those models. As we already explained, the twisted sector is independent on $e_{L,R}^i$ so the massless states for $e_{L,R}^i = 0$ remain massless for any value of $e_{L,R}^i \neq 0$. The $k=1$ model does not contain any massless twisted states, where as in the $k=0$ model

there are two massless chiral multiplets transforming as $(16, 1, 4, 1)$ under the $G_{k=0}$. So in total there are 8 $SO(10)$ families coming from the $k=0$ twisted sector and none from the $k=1$. In the untwisted sector, for both the $k=0$ and $k=1$ cases, the chiral multiplets, i.e. the scalar superfields coming from the truncation of $Spin_{16}$, get charges $\pm e$. For what we saw before, this implies that the $(-)^{b_s}$ projection selects a particular linear combination of states in the $(16 + 16, 4 + 4, 1, 1)$ of the $N=1$ broken gauge group. The physical states are, alternatively, classified in terms of representations of the unbroken gauge group: they belong to The $(Spin_8, Spin_8, 1, 1)$ of G_k . This representation is nonchiral, obviously, so they can get a mass, which is easily found to be $m = |e|$ (here we put $\sqrt{2}R = 1$).

The gauge supermultiplets corresponding to broken generators of the symmetry group acquire a mass $m = |e|$, together with the Higgs fields necessary to provide the degrees of freedom of a $N=1$ massive vector multiplet (1 massive spin 1, 2 spin $\frac{1}{2}$, 1 spin 0). The remaining light fields, with nonzero mass, are two scalar supermultiplets. They are in the fundamental of $SO(5) \times SO(21)$, in the $k=0$ case, and in the fundamental of $SO(13) \times SO(13)'$ for $k=1$. The light-mass field spectrum therefore, shows the characteristic features associated with spontaneous symmetry breaking. The particular example we have examined, albeit a simplistic one, is already interesting, because it corresponds to the breaking of a chiral gauge symmetry, a step which can not be avoided in any string model with some pretensions of being realistic.

In this paper we have presented a new mechanism which yields such a result.[15] The same mechanism can be applied to more general cases to break spontaneously gauge symmetries and, possibly, supersymmetry by introducing non zero e_i^j charge parameters.

All the fundamental requirements of the underlying superconformal two dimensional theory are satisfied in the proposed $T(x_s)/Z_2$ models; namely, absence of local anomalies multiloop modular invariance and factorization of the higher genus partition function on the boundaries of moduli space. The correct connection between spin statistics and the correct particle interpretation of the one loop partition function follow as a consequence.

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