Superspace Action for the First Massive States of the Superstring

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Using the manifestly spacetime-supersymmetric version of open superstring field theory, we construct the free action for the first massive states of the open superstring compactified to four dimensions. This action is in N=1 D=4 superspace and describes a massive spin-2 multiplet coupled to two massive scalar multiplets.

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1. Introduction

The conventional action for open superstring field theory contains contact-term problems caused by the presence of picture-changing operators\[1\]. Recently, a new action was proposed for open superstring field theory where these picture-changing operators are absent\[2\]. In addition to eliminating the contact-term problems, this new action can be written in a form which is manifestly N=1 D=4 spacetime-supersymmetric.

After compactifying the open superstring on a Calabi-Yau manifold to four dimensions, it is easy to show that the massless compactification-independent contribution to the free part of this action reproduces the usual N=1 D=4 superspace action for super-Maxwell. In this paper, the first massive contribution to the free part of this action will be computed in N=1 D=4 superspace.

In an earlier paper\[3\], we used open superstring vertex-operator arguments to show that the massive D=4 spin-two multiplet can be described in N=1 D=4 superspace by a vector superfield $V_m$ satisfying the constraint $D^m\sigma^m_{\alpha\dot{\alpha}}V_m = (\partial^m\partial_n + 2)V_m = 0$. It will be shown here that these constraints come from gauge-fixing the equations of motion of a superspace action involving not just $V_m$, but also a spinor and two scalar superfields, $V_\alpha$, $B$, and $C$. The propagating degrees of freedom of this action describe a massive spin-2 multiplet as well as two massive scalar multiplets which are always present in Calabi-Yau compactifications of the open superstring to four dimensions.

Before constructing the open superstring field theory action for these massive states, we shall review the construction of the action for the massless states.

2. Construction of the Free Action for the Massless States

The topological description of the superstring, developed by one of the authors with Vafa in reference \[4\], is particularly suitable for formulating open superstring field theory\[2\]. In this approach, the BRST current is replaced by a spin-one generator $G^+$, the $\eta$ ghost (which comes from bosonizing the RNS super-reparameterization ghosts as $\beta = \partial\xi e^{-\phi}$ and $\gamma = \eta e^\phi$) is replaced by a spin-one generator $\tilde{G}^+$, and the ghost number current is replaced by a spin-one generator $J$. As discussed in \[4\], these three generators form part of a ‘small’ twisted N=4 algebra.

The usual condition for physical vertex operators is $Q(V) = 0$ where $V$ is independent of the $\xi$ zero mode, i.e. $\eta(V) = 0$. Note that $G(V)$ always means the contour integral of $G$ around $V$. Since the $\eta$ cohomology is trivial, one can always find an operator $\Phi$ such
that \( V = \eta(\Phi) \). So the condition for \( \Phi \) to be physical is that \( \tilde{G}^+(G^+(\Phi)) = 0 \) where \( \Phi \) is defined up to the gauge invariance \( \delta \Phi = G^+(\Lambda) + \tilde{G}^+(\bar{\Lambda}) \). These linearized equations of motion and gauge invariances are easily obtained from the string field theory action

\[
S = \langle \Phi \ G^+(\tilde{G}^+(\Phi)) \rangle
\]  

(2.1)

where \( \langle \rangle \) is the two-point correlation function on a sphere. As usual in open superstrings, the states carry Chan-Paton factors which will be suppressed throughout the paper.

The manifestly spacetime-supersymmetric description of the superstring is related to the usual RNS description by a field redefinition. For compactifications of the superstring on a Calabi-Yau threefold, this field redefinition allows \( N=1 \) \( D=4 \) super-Poincaré invariance to be made manifest. The field redefinition takes the left-moving worldsheet matter and ghost fields of the RNS description into five free bosons \((x^m, \rho)\) (where \( m = 0 \) to 3) which satisfy the OPE’s

\[
x^m(y)x^n(z) \to \log|y - z|\eta^{mn}, \quad \rho(y)\rho(z) \to \log(y - z),
\]

(2.2)

eight free fermions \((\theta^\alpha, \bar{\theta}^{\dot{\alpha}}, p^\alpha, \bar{p}^{\dot{\alpha}})\) (where \( \alpha \) and \( \dot{\alpha} = 1 \) or 2) which satisfy the OPE’s

\[
p_\alpha(y)\theta^\beta(z) \to \delta_\alpha^\beta(y - z)^{-1}, \quad \bar{p}^{\dot{\alpha}}(y)\bar{\theta}^{\dot{\beta}}(z) \to \delta^{\dot{\alpha}}_{\dot{\beta}}(y - z)^{-1},
\]

and a field theory for the six-dimensional compactification manifold which is described by the \( c = 3 \) \( N=2 \) superconformal generators \([T_C, G^+_C, G^-_C, J_C]\). This \( N=2 \) superconformal field theory is twisted so that \( G^+_C \) has conformal weight +1 and \( G^-_C \) has conformal weight +2. As usual, right-moving fields are related to the left-moving fields through boundary conditions.

The ‘small’ \( N=4 \) superconformal generators are defined in terms of these free fields by

\[
T = \frac{1}{2} \partial x^m \partial x_m + p_\alpha \partial \theta^\alpha + \bar{p}^{\dot{\alpha}} \partial \bar{\theta}^{\dot{\alpha}} + \frac{1}{2} \partial \rho \partial \rho - \frac{i}{2} \partial^2 \rho + T_C
\]

\[
G^+ = e^{i\rho} d^\alpha d_\alpha + G^+_C, \quad G^- = e^{-i\rho} \bar{d}^{\dot{\alpha}} \bar{d}_{\dot{\alpha}} + G^-_C,
\]

(2.3)

\[
\tilde{G}^+ = e^{-2i\rho+iH_C} \bar{d}^{\dot{\alpha}} \bar{d}_{\dot{\alpha}} + e^{-i\rho+iH_C}(G^-_C),
\]

\[
\tilde{G}^- = e^{2i\rho-iH_C} d^\alpha d_\alpha + e^{i\rho-iH_C}(G^+_C),
\]

\[
J = i\partial \rho + iJ_C, \quad J^{++} = e^{-i\rho+iH_C}, \quad J^{--} = e^{i\rho-iH_C},
\]
where
\[
d_\alpha = p_\alpha + \frac{i}{2} \tilde{\theta}^\alpha \partial x_{\alpha \dot{\alpha}} - \frac{1}{4} (\tilde{\theta})^2 \partial \theta_\alpha + \frac{1}{8} \theta_\alpha \partial (\tilde{\theta})^2, \tag{2.4}
\]
\[
d_{\dot{\alpha}} = \bar{p}_{\dot{\alpha}} + \frac{i}{2} \theta^\alpha \partial x_{\alpha \dot{\alpha}} - \frac{1}{4} (\theta)^2 \partial \bar{\theta}_{\dot{\alpha}} + \frac{1}{8} \bar{\theta}_{\dot{\alpha}} \partial (\theta)^2,
\]
\(J_C = \partial H_C\), and we use the bispinor convention \(v_{\alpha \dot{\alpha}} = \sigma^m_{\alpha \dot{\alpha}} v_m\). Our Pauli matrices, \(\sigma^m_{\alpha \dot{\alpha}}\) and \(\tilde{\sigma}^m_{\dot{\alpha} \alpha}\), are those of Wess and Bagger\(^4\), but we choose to define \((\theta)^2 = (\theta^\alpha \theta_\alpha)\) and \((\tilde{\theta})^2 = (\tilde{\theta}^{\dot{\alpha}} \tilde{\theta}_{\dot{\alpha}})\), and to use the ‘mostly-negative’ Minkowski metric \(\eta^{mn} = \text{diag}(+1, -1, -1, -1)\) so that \(Tr(\sigma^m \tilde{\sigma}^n) = 2\eta^{mn}\).

Although our hermiticity conditions on the superspace variables are \((x^m)^* = x^m\) and \((\theta^\alpha)^* = \tilde{\theta}^{\dot{\alpha}}\) as usual, our hermiticity conditions on the chiral bosons are
\[
\rho^* = 2\rho - H_C, \quad H_C^* = 3\rho - 2H_C,
\]
so that \((G^+)^* = \bar{G}^+\). Note that \(d_\alpha\) and \(d_{\dot{\alpha}}\) have been defined to commute with the spacetime supersymmetry generators and to satisfy the OPE’s
\[
d_\alpha(y) d_{\dot{\alpha}}(z) \rightarrow i \frac{\Pi_{\alpha \dot{\alpha}}}{y - z}, \quad d_\alpha(y) \Pi_{\beta \dot{\beta}}(z) \rightarrow -2i \epsilon_{\alpha \beta} \partial \bar{\theta}_{\dot{\beta}} \frac{y - z}{y - z}, \tag{2.5}
\]
where
\[
\Pi_{\alpha \dot{\alpha}} = \partial x_{\alpha \dot{\alpha}} + i \theta_\alpha \partial \bar{\theta}_{\dot{\alpha}} + i \bar{\theta}^{\dot{\alpha}} \partial \theta_\alpha.
\]

To make N=1 D=4 supersymmetry manifest in the string field theory action, one integrates over the non-zero modes of the worldsheet fields in the correlation function of (2.7), but leaves the integration over the zero modes of \(x^m, \theta^\alpha, \) and \(\tilde{\theta}^{\dot{\alpha}}\) explicit in the action. The result is the N=1 D=4 superspace action
\[
S = \int d^4 x d^2 \theta d^2 \tilde{\theta} < \Phi G^+(\bar{G}^+(\Phi)) > \tag{2.6}
\]
where <> now does not include the contribution of the zero modes of \(x^m, \theta^\alpha, \) and \(\tilde{\theta}^{\dot{\alpha}}\).

As our first example, we shall consider the contribution to the above action from the massless compactification-independent states. Since there is no tachyon, the massless states at zero momentum are described by string fields of conformal weight zero. The string field is independent of the compactification manifold and is of neutral U(1)-charge (i.e. zero ghost number), so it is described by a scalar superfield \(V(x, \theta, \bar{\theta})\). Plugging \(\Phi = V\) into the action of (2.6), one reproduces the super-Maxwell action
\[
S = \int d^4 x d^2 \theta d^2 \tilde{\theta} V \bar{D}_{\dot{\alpha}} D^2 \bar{D}_{\dot{\alpha}} \tilde{V} \tag{2.7}
\]
where \(D_\alpha = \frac{\partial}{\partial \theta_\alpha} + \frac{i}{2} \tilde{\theta}^\alpha \sigma^m_{\alpha \dot{\alpha}} \partial_m\) and \(D_{\dot{\alpha}} = \frac{\partial}{\partial \bar{\theta}^{\dot{\alpha}}} + \frac{i}{2} \theta^\alpha \sigma^m_{\alpha \dot{\alpha}} \partial_m\). Furthermore, the gauge invariances \(\delta \Phi = G^+(\Lambda) + \bar{G}^+(\bar{\Lambda})\) reproduce the expected gauge invariances
\[
\delta V = D^2 \lambda + \bar{D}^2 \bar{\lambda}
\]
where \(\Lambda = e^{-i \rho} \lambda\) and \(\bar{\Lambda} = e^{2i \rho - i H_C} \bar{\lambda}\). 

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3. Action for First Massive States

The field theory action of (2.6) will now be used to compute the contribution of the first massive states of the open superstring. We shall consider a generic Calabi-Yau three-fold and will not consider states which depend on the explicit structure of the compactification manifold. However, we will allow states which can be constructed from the Calabi-Yau U(1) current $\partial H_C$. As will be shown later, the presence of such states in the action is necessary for gauge invariance.

Since we want to describe states of $(mass)^2 = 2$, our string field should contain conformal weight +1 at zero momentum. Furthermore, it should have no U(1) charge and should only depend on the compactification manifold through the U(1) current. The most general such string field is

$$\Phi = d^\alpha W_\alpha(x, \theta, \bar{\theta}) + \bar{d}_{\dot{\alpha}} \bar{W}_{\dot{\alpha}}(x, \theta, \bar{\theta}) + \Pi^m V_m(x, \theta, \bar{\theta}) + \partial \theta^\alpha V_\alpha(x, \theta, \bar{\theta}) + \partial \bar{\theta}^{\dot{\alpha}} \bar{V}_{\dot{\alpha}}(x, \theta, \bar{\theta})$$

$$+ i (\partial \rho - \partial H_C) B(x, \theta, \bar{\theta}) + (\partial H_C - 3 \partial \rho) C(x, \theta, \bar{\theta})$$

(3.1)

where $B$ and $C$ are real superfields.

This field transforms as $\delta \Phi = G^+(\Lambda) + \bar{G}^+ (\bar{\Lambda})$ under the gauge transformations parameterized by

$$\Lambda = e^{-i\rho} (d^\alpha C_\alpha + \bar{d}_{\dot{\alpha}} \bar{E}_{\dot{\alpha}} + \Pi^m B_m + \partial \rho F + \partial \theta^\alpha B_\alpha + \partial \bar{\theta}^{\dot{\alpha}} \bar{H}_{\dot{\alpha}}),$$

$$\bar{\Lambda} = e^{2i\rho - iH_C} (d^\alpha E_\alpha + \bar{d}_{\dot{\alpha}} \bar{C}_{\dot{\alpha}} + \Pi^m \bar{B}_m + (2 \partial \rho - \partial H_C) \bar{F} + \partial \theta^\alpha H_\alpha + \partial \bar{\theta}^{\dot{\alpha}} \bar{B}_{\dot{\alpha}}).$$

Note that $\delta \Phi = 0$ when $\Lambda = G^+ \Omega$ or $\bar{\Lambda} = \bar{G}^+ \bar{\Omega}$. Defining

$$\Omega = e^{-2i\rho} (\partial^2 \theta^\alpha M_\alpha + (\partial \theta)^2 N + \partial \theta^\alpha \bar{d}_{\dot{\alpha}} \tilde{P}_{\alpha \dot{\alpha}} + \Pi_{\alpha \dot{\alpha}} \partial \theta^\alpha \tilde{Q}_{\dot{\alpha}}),$$

one can see that the transformations parameterized by $C_\alpha$, $B_m$, $F$, and $H_\alpha$ can be ignored. Furthermore, the transformations parameterized by $B_\alpha$ and $\bar{B}_{\dot{\alpha}}$ can be used to algebraically gauge-fix $W_\alpha = \bar{W}_{\dot{\alpha}} = 0$ in the string field $\Phi$.

In this gauge, the remaining string fields transform under the remaining gauge transformations as:

$$\delta V_m = -2i (\sigma_m)_{\alpha \dot{\alpha}} D^\alpha \bar{E}_{\dot{\alpha}} - 2i (\sigma_m)_{\alpha \dot{\alpha}} \bar{D}_{\dot{\alpha}} E^\alpha,$$

(3.2)

$$\delta V_\alpha = 4E_\alpha - \frac{1}{2} D^2 \bar{D}^2 E_\alpha + i \partial_{\alpha \dot{\alpha}} D^2 \bar{E}_{\dot{\alpha}},$$

$$\delta (C + iB) = iD^\alpha \bar{D}^2 E_\alpha.$$
Note that $B$ and $C$ are not gauge-invariant, so they can not be dropped from the action without breaking gauge invariance. In our earlier analysis of massive vertex operators \[3\], it was not necessary to include $B$ and $C$ since we were working in a fixed gauge.

Finally, the string field action $S = \int d^4x d^2\theta d^2\bar{\theta} < \Phi G^+ (\bar{G}^+(\Phi)) >$ in the gauge $W_\alpha = \bar{W}_\alpha = 0$ is given by

\[
S = \int d^4x d^2\theta d^2\bar{\theta} [-V^m (\bar{D}^2 D^2 V_m - 4V_m - 2i(\bar{\sigma}^n)^{\hat{\alpha}} \partial_n \bar{D}_\hat{\alpha} D_\alpha V_m \tag{3.3}
-4i(\bar{\sigma}_m)^{\hat{\alpha}} (D_\alpha \bar{V}_\hat{\alpha} + D_{\hat{\alpha}} V_\alpha) - 8\delta_m B + 12(\bar{\sigma}_m)^{\hat{\alpha}} [D_{\hat{\alpha}}, D_\alpha] C
+ V^\alpha (-4\bar{D}^\hat{\alpha} D_\alpha \bar{V}_\hat{\alpha} - 2\bar{D}^2 V_\alpha + 24\partial_{\alpha\hat{\alpha}} \bar{D}^\hat{\alpha} C - D_\alpha \bar{D}^2 (18iC + 2B))
+ \bar{V}^\hat{\alpha} (-2D^2 \bar{V}_\hat{\alpha} - 24\delta_{\alpha\hat{\alpha}} D^\alpha C + \bar{D}_\hat{\alpha} D^2 (18iC - 2B))
- 3C (11\bar{D}^2 D^2 - 16\delta^m \partial_m - 8) C + B (-\bar{D}^2 D^2 - 8) B + 6B \delta_{\alpha\hat{\alpha}} [\bar{D}^\hat{\alpha}, D^\alpha] C ].
\]

To evaluate the propagating degrees of freedom, it is convenient to use the gauge transformations parameterized by $E_\alpha$ and $\bar{E}_{\hat{\alpha}}$ to gauge-fix $V_\alpha = \bar{V}_{\hat{\alpha}} = 0$. Note that this gauge-fixing involves derivatives, so unlike the gauge-fixing of $W_\alpha$, it cannot be performed directly in the action. Furthermore, $\delta V_\alpha = 0$ when the gauge parameter $E_\alpha$ satisfies $E_\alpha = \frac{i}{2} \partial_{\alpha\hat{\alpha}} D^2 E^\hat{\alpha}$ and $(\partial^n \partial_n - 1) E_\alpha = 0$. This implies that there is a residue gauge transformation even after gauge-fixing $V_\alpha = 0$.

In the gauge $V_\alpha = \bar{V}_{\hat{\alpha}} = 0$, the equations of motion from the action of \[3.3\] are

\[
\frac{1}{2} \{ \bar{D}^2, D^2 \} V_m - 4V_m - 2\delta^m \partial_n V_m - 4\delta_m B + 6(\bar{\sigma}_m)^{\hat{\alpha}} [\bar{D}_\hat{\alpha}, D_\alpha] C = 0, \tag{3.4}
4i(\sigma^m)_{\alpha\hat{\alpha}} \bar{D}^\hat{\alpha} V_m + 24\partial_{\alpha\hat{\alpha}} \bar{D}^\hat{\alpha} C + D_\alpha \bar{D}^2 (-18iC - 2B) = 0, \tag{3.5}
-8\delta^m V_m - 16B - \{ \bar{D}^2, D^2 \} B + 6\delta_{\alpha\hat{\alpha}} [\bar{D}_\hat{\alpha}, D^\alpha] C = 0, \tag{3.6}
-12\delta^m [\bar{D}_\hat{\alpha}, D_\alpha] V_m + 48C - 33\{ \bar{D}^2, D^2 \} C + 96\delta_m \partial^m C - 6\delta_{\alpha\hat{\alpha}} [\bar{D}^\hat{\alpha}, D^\alpha] B = 0. \tag{3.7}
\]

Plugging \[3.5\] into \[3.6\] and \[3.7\], one learns that

\[
B = \frac{i}{8} [D^2, \bar{D}^2] C \tag{3.8}
16(\partial^m \partial_m - 1) C = 6\{ D^2, \bar{D}^2 \} C. \tag{3.9}
\]

The equation of motion for $C$ has two solutions:

\[
a) \quad D^2 C = 0, \quad (\partial^m \partial_m - 1) C = 0; \tag{3.10}
\]

\[
b) \quad D^2 C = 0, \quad (\partial^m \partial_m - 1) C = 0; \tag{3.11}
\]

\[
c) \quad D^2 C = 0, \quad (\partial^m \partial_m - 1) C = 0; \tag{3.12}
\]

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Using the residue gauge transformations of $E_\alpha$, any solution of equation (3.10) can be gauged away. Therefore, the only physical degrees of freedom of $C$ come from solutions of (3.11), which are of the form:

$$C = F + \bar{F}$$

(3.12)

where $F$ is a chiral superfield satisfying

$$\bar{D}_\dot{\alpha}F = D_\alpha F = (\partial^m \partial_m + 2)F = (\partial^m \partial_m + 2)\bar{F} = 0.$$

Although $F$ is a single chiral superfield, it describes two massive scalar multiplets since it does not satisfy the on-shell equation $D^2 F = 0$. Expanding in components,

$$F = X + \theta^\alpha \xi_\alpha + (\theta)^2 Y$$

where $X$ and $Y$ are complex bosons satisfying $(\partial^m \partial_m + 2)X = (\partial^m \partial_m + 2)Y = 0$. Defining $\psi_\dot{\alpha} = -\frac{i}{\sqrt{2}}\partial_\alpha \xi^\alpha$, the equations of motion for the fermion are

$$\partial_\alpha \xi^\alpha = i\sqrt{2} \psi_\dot{\alpha}, \quad \partial_\alpha \psi_\dot{\alpha} = i\sqrt{2} \xi^\alpha$$

which describe a massive Dirac spinor $\Lambda^A = (\xi^\alpha, \psi_\dot{\alpha})$.

Plugging (3.12) and (3.8) into (3.4) and (3.5), one obtains

$$D^\alpha V_{\alpha\dot{\alpha}} - 2i\partial_{\alpha\dot{\alpha}}D^\alpha F = 0,$$

(3.13)

$$\frac{1}{2} \{ \bar{D}^2, D^2 \} V_m - 4V_m - 2\partial^n \partial_n V_m - 16i(F - \bar{F}) = 0.$$  

(3.14)

Finally, shifting $\dot{V}_m = V_m - 2i\partial_m (F - \bar{F})$, one can rewrite (3.13) and (3.14) as

$$D^\alpha \dot{V}_{\alpha\dot{\alpha}} = 0, \quad (\partial^m \partial_m + 2)\dot{V}_m = 0,$$

(3.15)

which were shown in [3] to be the gauge-fixed equations of motion of a massive spin-two multiplet. Therefore, the propagating degrees of freedom described by the action of (3.3) are a massive spin-two multiplet and two massive scalar multiplets.

On-shell, these multiplets contain 12 bosonic and 12 fermionic degrees of freedom. The bosonic degrees of freedom can be identified with the following 12 Neveu-Schwarz light-cone gauge vertex operators:

$$b^m_{\frac{m}{2}} |0\rangle, \quad b^m_{\frac{m}{2}} a^n_{-1} |0\rangle,$$

(3.16)
\[ \omega_{JKL} b^{-J}_{-\frac{1}{2}} b^{K}_{-\frac{1}{2}} b^{L}_{-\frac{1}{2}} |0\rangle, \quad \omega_{\bar{J}\bar{K}\bar{L}} b^{-\bar{J}}_{-\frac{1}{2}} b^{\bar{K}}_{-\frac{1}{2}} b^{\bar{L}}_{-\frac{1}{2}} |0\rangle, \]
\[ b^{m}_{-\frac{1}{2}} g_{J\bar{J}} b^{J}_{-\frac{1}{2}} b^{\bar{J}}_{-\frac{1}{2}} |0\rangle, \quad g_{J\bar{J}} b^{J}_{-\frac{1}{2}} a^{\bar{J}}_{-\frac{1}{2}} |0\rangle, \quad g_{J\bar{J}} b^{J}_{-\frac{1}{2}} a^{\bar{J}}_{-\frac{1}{2}} |0\rangle, \]

where \( m, n, p = 2, 3 \) are the D=4 light-cone indices, \( g_{J\bar{J}} \) and \( \omega_{JKL} \) are the metric and holomorphic 3-form used to define the Calabi-Yau manifold, and \( b^{m}_{N} \) and \( a^{m}_{N} \) are the oscillator modes of the light-cone \( \psi^{m} \) and \( \partial X^{m} \). Similarly, the fermionic degrees of freedom can be identified with the following 12 Ramond light-cone gauge vertex operators:

\[ b^{m}_{-1} |0\rangle^{++++}, \quad b^{m}_{-1} |0\rangle^{------}, \quad a^{m}_{-1} |0\rangle^{-+++}, \quad a^{m}_{-1} |0\rangle^{++++}, \quad (3.17) \]
\[ g_{J\bar{J}} b^{J}_{-1} b^{\bar{J}}_{0} |0\rangle^{++-}, \quad g_{J\bar{J}} a^{\bar{J}}_{-1} b^{J}_{0} |0\rangle^{--+}, \]
\[ g_{J\bar{J}} b^{J}_{0} b^{\bar{J}}_{-1} |0\rangle^{-++}, \quad g_{J\bar{J}} b^{J}_{0} a^{\bar{J}}_{-1} |0\rangle^{++-} \]

where \( |0\rangle^{\pm\pm\pm\pm} \) is the light-cone spinor in SU(4) notation where the last 3 \( \pm \) signs refer to the six Calabi-Yau directions. Although there might be other states of \((mass)^2 = 2\) in the open superstring spectrum which come from Calabi-Yau excitations, the 12+12 states described above are always present in any Calabi-Yau compactification.

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References

[1] E. Witten, Nucl. Phys B276(1986) 291; C. Wendt, Nucl. Phys B314(1989) 209.

[2] N. Berkovits, Nucl. Phys B450(1995) 90, hep-th/9503093.

[3] N. Berkovits and M. M. Leite, Phys.Lett.B415(1997) 144, hep-th/9709148.

[4] N. Berkovits and C. Vafa, Nucl. Phys B433(1995) 123; see also N. Berkovits, “A New Description of the Superstring”, proceedings of the VIII J. A. Swieca Summer School on Particles and Fields, World Scientific Publishing (1996), hep-th/9604123; Nucl. Phys. B431(1994) 258, hep-th/9404162.

[5] J. Wess and J. Bagger, Supersymmetry and Supergravity: Notes from Lectures given at Princeton University, Princeton Univ. Press, 1982