

Note on Mediation of Supersymmetry Breaking from Closed to Open Strings

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Abstract

We discuss the mediation of supersymmetry breaking from closed to open strings, extending and improving previous analysis of the authors in Nucl. Phys. B 695 (2004) 103 [hep-th/0403293]. In the general case, we find the absence of anomaly mediation around any perturbative string vacuum. When supersymmetry is broken by Scherk-Schwarz boundary conditions along a compactification interval perpendicular to a stack of D-branes, the gaugino acquires a mass at two loops that behaves as $m_{1/2} \sim g^4 m_{3/2}^3$ in string units, where $m_{3/2}$ is the gravitino mass and g is the gauge coupling.

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The purpose of this note is to extend the previous discussion [1] of mediation of supersymmetry breaking between closed and open string sectors. At the same time, we will also improve and clarify some computations presented in [1].

The mediation of supersymmetry breaking is (by definition) the mechanism responsible for communicating supersymmetry breaking from a hidden sector to the spectrum of observable particles. Although there are several possible ways how such a mechanism may be realized in low-energy effective field theory, most of its concrete implementations involve non-renormalizable interactions, therefore the outcome can be sensitive to the ultraviolet completion of the theory. Hence the finite superstring theory offers a valuable framework for studying the effects of supersymmetry breaking while keeping the ultra-violet physics under control.

One type of mediation possible in this context, is the so-called anomaly mediation [2], providing a contribution to the gaugino mass that scales linearly with the gravitino mass: $m_{1/2} \sim b_0 g^2 m_{3/2}$, where g is the gauge coupling and b_0 is the coefficient of the corresponding one-loop beta-function. However, as explained in [1], this contribution is absent in any perturbative string vacuum. The reason is that such a result should arise at one-loop level, as dictated by the power of the gauge coupling, *e.g.* on a world-sheet with two boundaries (annulus) or one boundary and a crosscap (Möbius strip). The corresponding string diagram involving two left-handed gauginos at zero momentum vanishes though, due to the $U(1)$ charge conservation of the two-dimensional (2d) $N = 2$ superconformal symmetry. Indeed, the massless gaugino vertex operator of definite chirality α , at zero momentum, in the canonical $-1/2$ -ghost picture, reads:

$$V_\alpha^{(-1/2)}(x) =: e^{-\varphi/2} S_\alpha e^{i\frac{\sqrt{3}}{2}H} :, \quad (1)$$

where x is a position on the boundary of the world-sheet, φ is the scalar bosonizing the superghost system, H is the free 2d boson associated to the $N = 2$ $U(1)$ current $J = i\sqrt{3}\partial H$, and we neglected the Chan-Paton gauge indices for simplicity. The two-point function involves, besides the two gauginos of the same chirality at the boundary of the world-sheet (annulus or Möbius strip), one picture changing operator (PCO). The latter can provide at most -1 charge which is not sufficient for cancelling the $+3$ $U(1)$ charge of the gauginos and thus, the amplitude vanishes. Charge cancellation can be achieved at higher order, requiring a Riemann surface of Euler characteristic at least equal to -1 . An example of such a surface contains one handle and one boundary and will be studied in the example described below.

The particular setup considered in [1] is the gravitational mediation from the closed string sector with supersymmetry broken by Scherk-Schwarz [3] boundary conditions in one of the compact directions, which plays the role of the hidden sector, to the “observable” sector of open strings ending on D-branes perpendicular to the Scherk-Schwarz direction [4]. The corresponding gauginos that remain massless at the tree level, acquire masses due to the world-sheet diagram with one handle and one boundary – genus $g = 1$ Riemann surface Σ with $h = 1$ boundary, *i.e.* with Euler characteristic -1 , see Fig.1.

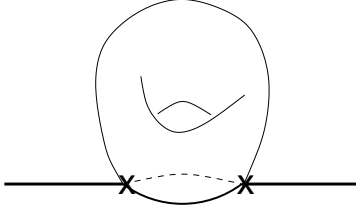


Figure 1: *Bordered $g = 1$ surface Σ with the two gaugino vertices inserted at the boundary*

In Ref.[1], we discussed type II compactifications on $T^2 \times K3$, with the Scherk-Schwarz circle of radius R embedded in T^2 . Although the generic case can be studied to some extent, the mass computations simplify enormously in the orbifold limit, with $K3$ represented as the quotients T^4/\mathbb{Z}_2 or T^4/\mathbb{Z}_N . Unfortunately, in this case the mass is protected by the orbifold symmetries – the remnants of the continuous internal rotations, or equivalently R -symmetries of the low-energy effective field theory – hence the result is zero. We will first review the origin of this result and then we will try to circumvent it by blowing up the orbifold singularity of $K3$.

The gaugino mass is given by the following integral [1]:

$$m_{1/2} = g_s^2 \int_{F(\Sigma)} d\mu \int_{\partial\Sigma} dx dy \mathcal{A}(x, y), \quad (2)$$

where $d\mu$ is the measure of the moduli space of Σ , with the integration extending over the fundamental domain $F(\Sigma)$, while the two-point function

$$\begin{aligned} \mathcal{A}(x, y) = \theta \begin{bmatrix} 0 \\ \frac{1}{2}\vec{1} \end{bmatrix} (x - \Delta) \frac{\sigma(x)\sigma(y)}{\prod_{I<J}^{3,4,5} E(z_I, z_J) \prod_{I=3}^5 \sigma^2(z_I)} \\ \times \prod_{I=3}^5 \theta_{h_I^{-1}} \begin{bmatrix} 0 \\ \frac{1}{2}\vec{1} \end{bmatrix} (z_I - \Delta) \partial X_{h_I}(z_I) \mathcal{Z}. \end{aligned} \quad (3)$$

is integrated over the boundary. In the mass formula (2), $g_s = g^2$ is the string coupling. The additional points, z_I , labeled by the internal planes $I = 3, 4, 5$, take values in the set of

the insertion points, z_a , $a = 1, 2, 3$, of PCOs, and the above expression should be summed over all permutations $\{z_{I(a)}\}$. Although *a priori* arbitrary, the PCO insertion points are subject to the constraint

$$\sum_{I=3}^{I=5} z_I = y + 2\Delta. \quad (4)$$

as a result of a gauge choice made in arriving to (3). Here, Δ is the Riemann θ -constant. Although the above gauge choice is formally not allowed [5], it can be used throughout the computations. Indeed, by inserting another vertex of an open string Wilson line, one can use an appropriate gauge condition and show that the amplitude can be written as the variation with respect to the Wilson line of the original one, evaluated by formally using the condition (4).

After summing over all 6 permutations, the two-point function (3) should become manifestly independent of the PCO insertion points. In Eq.(3), ∂X_{h_I} are the zero modes (instanton contributions) twisted by the orbifold group elements h_I while the position-independent factor \mathcal{Z} includes the lattice partition function as well as all non-zero mode determinants. Finally, σ is the one-differential with no zeroes or poles and E is the prime form. The crucial property of the prime form is the antisymmetry $E(z_I, z_J) = -E(z_J, z_I)$. As a result, the permutation sum amounts to antisymmetrizing the factor

$$\theta_{h_3}^{-1} \begin{bmatrix} 0 \\ \frac{1}{2} \vec{1} \end{bmatrix} (z_3 - \Delta) \partial X_{h_3}(z_3) K(z_4, z_5), \quad K(z_4, z_5) = \prod_{I=4}^5 \theta_{h_I}^{-1} \begin{bmatrix} 0 \\ \frac{1}{2} \vec{1} \end{bmatrix} (z_I - \Delta) \partial X_{h_I}(z_I) \quad (5)$$

in the positions z_a . Note that, for the $T^2 \times K_3$ compactification under consideration, $h_3 = 1$ while $h_4 = h$, $h_5 = h^{-1}$, where h is the element of the $K3$ orbifold group \mathbb{Z}_N . Clearly, in the case of \mathbb{Z}_2 , $h = h^{-1}$, hence the result vanishes upon antisymmetrization. As shown in [1], similar conclusion holds for arbitrary \mathbb{Z}_N , at least up to terms that are exponentially suppressed in the large Scherk-Schwarz radius limit. As announced before, we will try to avoid this conclusion by blowing up the orbifold singularity. This can be achieved by switching on the vacuum expectation value of one of the blowing-up modes. Thus the amplitude will now include an additional insertion of the vertex operator creating the twisted blowing-up mode \mathcal{B} at zero momentum. Note that this additional vertex must be inserted in the 0-ghost picture in order to preserve the balance of the background ghost charge.

In the -1 -ghost picture, the zero momentum vertex of a blowing-up mode \mathcal{B} associated to the twisted sector (h, h^{-1}) , with $h = e^{2i\pi\epsilon}$ and $\epsilon = k/N$, reads:

$$V_{\mathcal{B}}^{(-1,-1)}(\zeta, \bar{\zeta}) =: e^{-i\epsilon H_4} e^{-i(1-\epsilon)\tilde{H}_4} \sigma_{--}^{4,1-\epsilon} e^{-i(1-\epsilon)H_5} e^{-i\epsilon\tilde{H}_5} \sigma_{--}^{5,\epsilon} :, \quad (6)$$

where σ_{--} is the corresponding twist field of conformal dimension $\epsilon(1-\epsilon)/2$ in both left and right movers. Here, (H_4, H_5) and $(\tilde{H}_4, \tilde{H}_5)$ are the scalars bosonizing the left- and right-moving fermionic coordinates of $K3$, respectively.¹ In order to change the picture, we use the supercurrents

$$T_L = \sum_{I=4,5} (\partial X^I e^{-iH_I} + \partial \bar{X}^I e^{iH_I}) \quad T_R = \sum_{I=4,5} (\bar{\partial} X^I e^{-i\tilde{H}_I} + \bar{\partial} \bar{X}^I e^{i\tilde{H}_I}). \quad (7)$$

We use the OPE rules [6]

$$\begin{aligned} \sigma_{--}^{I,\epsilon}(z, \bar{z}) \partial \bar{X}_I(\bar{w}) &\sim (z-w)^{-\epsilon} \sigma_{+-}^{I,\epsilon}(z, \bar{z}) & \sigma_{--}^{I,\epsilon}(z, \bar{z}) \bar{\partial} \bar{X}_I(\bar{w}) &\sim (\bar{z}-\bar{w})^{-1+\epsilon} \sigma_{-+}^{I,\epsilon}(z, \bar{z}) \\ \sigma_{--}^{I,\epsilon}(z, \bar{z}) \partial X_I(w) &\sim (z-w)^{-1+\epsilon} \sigma_{+-}^{I,\epsilon}(z, \bar{z}) & \sigma_{--}^{I,\epsilon}(z, \bar{z}) \bar{\partial} X_I(w) &\sim (\bar{z}-\bar{w})^{-\epsilon} \sigma_{-+}^{I,\epsilon}(z, \bar{z}) \end{aligned} \quad (8)$$

with the further short-distance expansion

$$\begin{aligned} \sigma_{-+}^{I,\epsilon}(z, \bar{z}) \partial \bar{X}_I(\bar{w}) &\sim (z-w)^{-\epsilon} \sigma_{++}^{I,\epsilon}(z, \bar{z}) & \sigma_{-+}^{I,\epsilon}(z, \bar{z}) \bar{\partial} \bar{X}_I(\bar{w}) &\sim (\bar{z}-\bar{w})^{-1+\epsilon} \sigma_{++}^{I,\epsilon}(z, \bar{z}) \\ \sigma_{-+}^{I,\epsilon}(z, \bar{z}) \partial X_I(w) &\sim (z-w)^{-1+\epsilon} \sigma_{++}^{I,\epsilon}(z, \bar{z}) & \sigma_{-+}^{I,\epsilon}(z, \bar{z}) \bar{\partial} X_I(w) &\sim (\bar{z}-\bar{w})^{-\epsilon} \sigma_{++}^{I,\epsilon}(z, \bar{z}) \end{aligned} \quad (9)$$

Using the $N = 2$ world-sheet supercurrent, one finds the blowing-up vertex operator in the 0-ghost picture:

$$\begin{aligned} V_{\mathcal{B}}^{(0,0)}(\zeta, \bar{\zeta}) =: & \sigma_{+-}^{4,1-\epsilon} e^{i(1-\epsilon)H_4} e^{-i(1-\epsilon)\tilde{H}_4} \sigma_{-+}^{5,\epsilon} e^{-i(1-\epsilon)H_5} e^{i(1-\epsilon)\tilde{H}_5} \\ & + \sigma_{-+}^{4,1-\epsilon} e^{-i\epsilon H_4} e^{i\epsilon\tilde{H}_4} \sigma_{+-}^{5,\epsilon} e^{i\epsilon H_5} e^{-i\epsilon\tilde{H}_5} \\ & + \sigma_{++}^{4,1-\epsilon} e^{i(1-\epsilon)H_4} e^{i\epsilon\tilde{H}_4} \sigma_{--}^{5,\epsilon} e^{-i(1-\epsilon)H_5} e^{-i\epsilon\tilde{H}_5} \\ & + \sigma_{--}^{4,1-\epsilon} e^{-i\epsilon H_4} e^{-i(1-\epsilon)\tilde{H}_4} \sigma_{++}^{5,\epsilon} e^{i\epsilon H_5} e^{i(1-\epsilon)\tilde{H}_5} : \end{aligned} \quad (10)$$

The modification of the amplitude due to the the insertion of the blowing-up vertex operator $\int d^2\zeta V_{\mathcal{B}}^{(0,0)}(\zeta, \bar{\zeta})$ can be obtained by repeating step by step the derivation of (3)

¹ Since the following paragraph corrects several (accumulative) misprints contained in [1], we deliberately change the notation and repeat the computation of the vertex operator from the scratch.

presented in [1]. The only change is in the $K3$ part of the amplitude where, for a given spin structure s , the following correlators appear:

$$\begin{aligned}
& \left\langle \sigma_{+-}^{4,1-\epsilon}(\zeta, \bar{\zeta}) \partial X^4(z_4) \right\rangle \left\langle e^{iH_4/2}(x) e^{iH_4/2}(y) e^{-iH_4}(z_4) e^{i(1-\epsilon)H_4}(\zeta) e^{-i(1-\epsilon)\tilde{H}_4}(\bar{\zeta}) \right\rangle_s \times \\
& \left\langle \sigma_{-+}^{5,\epsilon}(\zeta, \bar{\zeta}) \partial X^5(z_5) \right\rangle \left\langle e^{iH_5/2}(x) e^{iH_5/2}(y) e^{-iH_5}(z_5) e^{-i(1-\epsilon)H_5}(\zeta) e^{i(1-\epsilon)\tilde{H}_5}(\bar{\zeta}) \right\rangle_s \\
& + \left\langle \sigma_{-+}^{4,1-\epsilon}(\zeta, \bar{\zeta}) \partial X^4(z_4) \right\rangle \left\langle e^{iH_4/2}(x) e^{iH_4/2}(y) e^{-iH_4}(z_4) e^{-i\epsilon H_4}(\zeta) e^{i\epsilon \tilde{H}_4}(\bar{\zeta}) \right\rangle_s \times \\
& \left\langle \sigma_{+-}^{5,\epsilon}(\zeta, \bar{\zeta}) \partial X^5(z_5) \right\rangle \left\langle e^{iH_5/2}(x) e^{iH_5/2}(y) e^{-iH_5}(z_5) e^{i\epsilon H_5}(\zeta) e^{-i\epsilon \tilde{H}_5}(\bar{\zeta}) \right\rangle_s \\
& \sim \theta_{s,h_4} \left(\frac{1}{2}(x+y) - z_4 + (1-\epsilon)(\zeta - \bar{\zeta}) \right) \theta_{s,h_5} \left(\frac{1}{2}(x+y) - z_5 - (1-\epsilon)(\zeta - \bar{\zeta}) \right) \times \\
& \left[\frac{E(z_4, \bar{\zeta}) E(z_5, \zeta)}{E(z_4, \zeta) E(z_5, \bar{\zeta})} \right]^{1-\epsilon} \frac{1}{E(\zeta, \bar{\zeta})^{2(1-\epsilon)^2}} \left\langle \sigma_{+-}^{4,1-\epsilon}(\zeta, \bar{\zeta}) \partial X^4(z_4) \right\rangle \left\langle \sigma_{-+}^{5,\epsilon}(\zeta, \bar{\zeta}) \partial X^5(z_5) \right\rangle \\
& + \theta_{s,h_4} \left(\frac{1}{2}(x+y) - z_4 - \epsilon(\zeta - \bar{\zeta}) \right) \theta_{s,h_5} \left(\frac{1}{2}(x+y) - z_5 + \epsilon(\zeta - \bar{\zeta}) \right) \times \\
& \left[\frac{E(z_4, \zeta) E(z_5, \bar{\zeta})}{E(z_4, \bar{\zeta}) E(z_5, \zeta)} \right]^\epsilon \frac{1}{E(\zeta, \bar{\zeta})^{2\epsilon^2}} \left\langle \sigma_{-+}^{4,1-\epsilon}(\zeta, \bar{\zeta}) \partial X^4(z_4) \right\rangle \left\langle \sigma_{+-}^{5,\epsilon}(\zeta, \bar{\zeta}) \partial X^5(z_5) \right\rangle,
\end{aligned} \tag{11}$$

Note that due to the $H_{4,5}$ internal charge conservation, there are no correlators involving σ_{++} twist fields. The spin structure sum can be performed using the same gauge condition (4), with the result that the factor K of Eq.(5) is replaced by:

$$\begin{aligned}
K(z_4, z_5) & \rightarrow \int d^2\zeta K_\epsilon(z_4, z_5, \zeta, \bar{\zeta}) = \\
& \int d^2\zeta \theta_{h_4^{-1}} \begin{bmatrix} 0 \\ \frac{1}{2}\vec{1} \end{bmatrix} (z_4 - \epsilon(\zeta - \bar{\zeta}) - \Delta) \theta_{h_5^{-1}} \begin{bmatrix} 0 \\ \frac{1}{2}\vec{1} \end{bmatrix} (z_5 + \epsilon(\zeta - \bar{\zeta}) - \Delta) \\
& \times \left[\frac{E(z_4, \zeta) E(z_5, \bar{\zeta})}{E(z_4, \bar{\zeta}) E(z_5, \zeta)} \right]^\epsilon \frac{1}{E(\zeta, \bar{\zeta})^{2\epsilon^2}} \left\langle \sigma_{-+}^{4,1-\epsilon}(\zeta, \bar{\zeta}) \partial X^4(z_4) \right\rangle \left\langle \sigma_{+-}^{5,\epsilon}(\zeta, \bar{\zeta}) \partial X^5(z_5) \right\rangle \\
& + (4 \leftrightarrow 5, \epsilon \leftrightarrow 1 - \epsilon),
\end{aligned} \tag{12}$$

The above expression is no longer symmetric in $z_4 \leftrightarrow z_5$ (except for $\epsilon = 1/2 \mathbb{Z}_2$ twist) therefore it can survive the antisymmetrization. Next, we will estimate the magnitudes of the antisymmetric part and of the corresponding gaugino mass term.

We are interested in the limit of large Scherk-Schwarz radius R , *i.e.* the limit of low gravitino mass $m_{3/2} \sim 1/R$ in string units. As explained in [1], in the $R \rightarrow \infty$ limit, the dominant contribution to the gaugino mass comes from the $\tau_2 \rightarrow \infty$ region of the moduli

space describing the handle degeneration limit. In this limit, Σ degenerates into a disk with two punctures left-over from the handle. As usual, the disk can be mapped into the upper half of the complex plane. Then the distance between the punctures is controlled by the remaining (real) modulus l of Σ . The gaugino mass has the form [1]:

$$m_{1/2} \sim g^4 \int d\tau_2 \int \frac{dl}{(e^{2\pi l} - 1)} \frac{\Gamma}{(\tau_2 - l)^{3/2}} \sum_m \frac{mR^2}{(\tau_2 + l)^2} \exp\left(-\frac{m^2\pi R^2}{\tau_2 - l}\right). \quad (13)$$

Here, the sum is over the winding modes on the Scherk-Schwarz circle, and the factor $(\tau_2 - l)^{-3/2}$ includes $(\tau_2 - l)^{-1}$ from the corresponding zero modes and $(\tau_2 - l)^{-1/2}$ from the partition function. The factor $(\tau_2 + l)^{-2}$ originates from the the non-compact zero modes (four-dimensional momenta). The factor $(e^{2\pi l} - 1)^{-1}$ is the combined effect of the integration measure and of non-zero mode determinants. We denote by Γ any additional moduli-dependence that may appear as a result of antisymmetrizing $K_\epsilon(z_4, z_5, \zeta, \bar{\zeta})$, Eq.(12). Note that the integral over the modulus l is dominated by the $l \rightarrow 0$ region. In fact, if Γ does not vanish in this limit, the logarithmic divergence may give rise to the additional τ_2 dependence due to the cutoff $l > e^{-\pi\tau_2}$ [1]. Thus the key question is the $\tau_2 \rightarrow \infty, l \rightarrow 0$ behavior of Γ – its “double degeneration limit” *i.e.* the limit of two coalescing punctures on the disk.

In order to extract the leading $l \rightarrow 0$ behavior of $K_\epsilon(z_4, z_5, \zeta, \bar{\zeta})$, we can set $\zeta = \bar{\zeta} = 0$ inside the arguments of the theta functions in Eq.(12). Furthermore, the twist correlators are evaluated on the disk, and they are completely determined by the $SL(2, R)$ covariance:

$$\langle \sigma_{+-}^\epsilon(\zeta, \bar{\zeta}) \partial X(z) \rangle = \frac{(\zeta - \bar{\zeta})^{\epsilon^2}}{(z - \zeta)^2} \left(\frac{z - \zeta}{z - \bar{\zeta}} \right)^\epsilon, \quad \langle \sigma_{-+}^{(1-\epsilon)}(\zeta, \bar{\zeta}) \partial X(z) \rangle = \frac{(\zeta - \bar{\zeta})^{\epsilon^2}}{(z - \bar{\zeta})^2} \left(\frac{z - \bar{\zeta}}{z - \zeta} \right)^\epsilon \quad (14)$$

After taking the corresponding limit of the prime-forms, $E(w_1, w_2) \rightarrow (w_1 - w_2)^{-1}$, we obtain

$$K_\epsilon(z_4, z_5, \zeta, \bar{\zeta}) \rightarrow \theta_{h_4^{-1}} \begin{bmatrix} 0 \\ \frac{1}{2}\mathbb{1} \end{bmatrix} (z_4 - \Delta) \theta_{h_5^{-1}} \begin{bmatrix} 0 \\ \frac{1}{2}\mathbb{1} \end{bmatrix} (z_5 - \Delta) \frac{1}{(z_4 - \bar{\zeta})^2 (z_5 - \zeta)^2} + (4 \leftrightarrow 5) \quad (15)$$

Note that the ϵ -dependence has disappeared in this limit. The above function vanishes upon antisymmetrization, thus $\Gamma = \mathcal{O}(l)$ and the l -integral in (13) converges at $l = 0$. The dominant $\tau_2 \rightarrow \infty$ region yields

$$m_{1/2} \sim g^4 \int \frac{d\tau_2}{\tau_2^{7/2}} \sum_m mR^2 \exp\left(-\frac{m^2\pi R^2}{\tau_2}\right). \quad (16)$$

After rescaling $\tau_2 \rightarrow \tau_2 R^2$, the above expression yields²

$$m_{1/2} \sim \frac{g^4}{R^3} \sim g^4 m_{3/2}^3. \quad (17)$$

The mass (17) can be understood within the effective field theory by looking at a generic one-loop graph involving a gravitational exchange. Each vertex brings one power of the Planck mass M_P in the denominator and is quadratically divergent in the ultraviolet, thus $m_{1/2} \sim m_{3/2} \Lambda_{UV}^2 / M_P^2$, where Λ_{UV} is the ultraviolet cutoff. This cutoff should be of order of the supersymmetry breaking scale [7],³ $\Lambda_{UV} \sim m_{3/2}$, hence $m_{1/2} \sim m_{3/2}^3 / M_P^2$ [8]. The result (17) confirms this expectation.

To summarize, the mediation of supersymmetry breaking from closed to open string sectors provides a superstring realization of the so-called gravitational mediation [9]. Another type of mediation, the anomaly mediation [2] discussed in the beginning and in [1], is absent, at least at the leading $\mathcal{O}(g_s) \sim \mathcal{O}(g^2)$ order. Moreover, it is also absent at the $\mathcal{O}(g^4)$ order, as follows from the result (17). A new type of non-gravitational mediation between open string sectors has been recently discussed in [5].

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² This result corrects [1], where a different conclusion has been reached without appropriate analysis of the twisted correlators.

³ We are grateful to Savas Dimopoulos and to Mary K. Gaillard for insisting on this point.

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