Scattering, Sequestering and Sort-Of Sequestering of Local Moduli in Local Models

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Based on arXiv:1109.xxxx JC, L. Witkowski

Acknowledgements

Talk is based on arXiv:1109.xxxx (together with Lukas Witkowski),

See parallel talk by LW on Tuesday 3.50pm for more details.

Related work: JC 0710.9873, Choi Jeong Okomura 0804.4283, Acharya Bobkov 0810.3285, Kane Kumar Shao 0905.2986, Blumenhagen JC Krippendof Moster Quevedo 0906.3297, JC Pedro 1003.0388, Choi Nilles Shin Trapletti 1011.0999, Berg Marsh McAllister Pajer 1012.1858

Talk Structure

- Motivation: Supersymmetry Breaking and Sequestering
- Quantum Correlator
- Classical Correlator and Sort-Of Sequestering

Talk

This talk studies disk correlators of the form

$$\langle \tau_s^{(1)} \tau_s^{(2)} \ldots \tau_s^{(n)} \psi\psi\phi \rangle$$

where $\psi\psi\phi$ is an open string Yukawa coupling on a D3 brane (at a singularity) and $\tau_s^{(i)}$ is a closed string blow-up mode.

τ_s may correspond to a mode located either at or far distant from the D3 singularity.

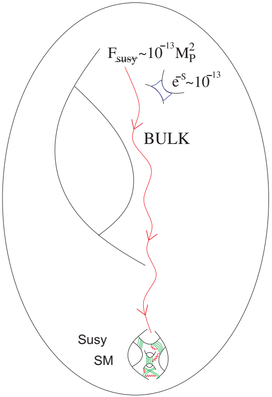
For $n = 1$ we calculate the full quantum and classical correlator.

For $n \geq 1$ we obtain 'sort-of sequestering' for distant moduli.

We motivate the calculation, state the quantum correlator and explain sequestering for distant moduli.

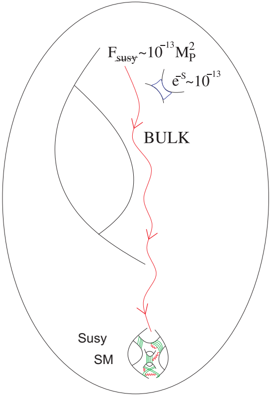
Motivation

Large Volume Scenario: MSSM realised locally on supersymmetric singularity, supersymmetry broken in bulk and on distant cycles.



Motivation

Soft terms come from communication from susy-breaking bulk/small moduli to MSSM sector:



Motivation

Dilaton and complex structure moduli are flux-stabilised in a supersymmetric fashion,

$$D_S W = 0, \qquad D_U W = 0.$$

Kähler moduli break supersymmetry through the structure of the LVS supergravity potential.

Structure of susy breaking is dominantly (up to $\mathcal{O}(\mathcal{V}^{-1})$) no-scale structure:

$$\begin{aligned} K &= -2 \ln \mathcal{V} \qquad (\sim -3 \ln(T + \bar{T})), \\ W &= W_0. \end{aligned}$$

Leading F-term is overall volume modulus,

$$|F^b|^2 = 3 m_{3/2}^2 \left(1 + \mathcal{O}(\mathcal{V}^{-\lambda}) + \ldots\right).$$

Motivation

Original analyses of soft terms in LVS considered matter metric

$$Z_{\alpha\beta} \sim \frac{1}{\mathcal{V}^{\mu}}.$$

For $\mu \neq 2/3$, soft terms are $\mathcal{O}(m_{3/2})$. However locality of physical Yukawa couplings

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_{\alpha}\tilde{K}_{\beta}\tilde{K}_{\gamma})^{\frac{1}{2}}}.$$

requires

$$Z \sim \frac{1}{\mathcal{V}^{2/3}}.$$

No-scale susy breaking: $\mathcal{O}(m_{3/2})$ tree-level soft terms cancel and it is necessary to analyse terms subleading in α' and g_s.

Motivation

Full sequestering does not occur in string compactifications.

However 'sort-of sequestering' occurs if

$$Z \sim e^{\hat{K}(T,\bar{T})/3}.$$

when the soft terms vanish.

Physical Yukawas behave as

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{\sqrt{Z_\alpha Z_\beta Z_\gamma}}$$

'Sort-of sequestering' of a distant Kähler modulus τ_s is equivalent to the independence of $\hat{Y}_{\alpha\beta\gamma}$ from τ_s.

Soft terms vanish to the order at which 'sort-of sequestering' holds.

What we do

By computing correlators

$$\langle \tau_s^{(1)} \tau_s^{(2)} ... \tau_s^{(n)} \psi \psi \phi \rangle$$

we directly probe the dependence of physical Yukawa couplings on blow-up moduli τ_s.

This is a direct test of 'sort-of sequestering' for these moduli.

By working in disk CFT we obtain results valid to all orders in α' and leading order in g_s.

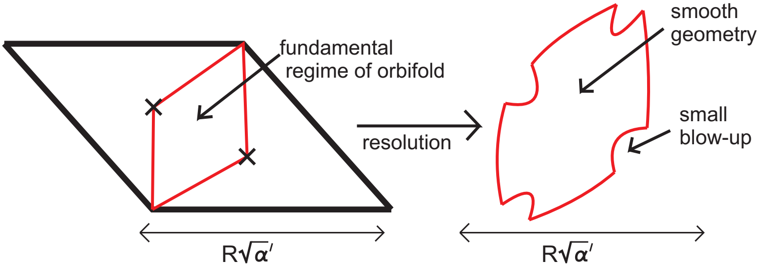
By including multiple powers of τ_s we also probe the limit where τ_s obtains a vev and resolves the orbifold to a smooth space.

Blow-up moduli τ_s can be located

- At the same singularity as MSSM fields.
- At geographically separated singularities.

What we do

Bonus: resolved toroidal orbifolds capture the same Swiss-cheese geometry that appears in the Large Volume Scenario:



Quantum Correlator

Consider the scattering of 3 open string (two fermionic, one bosonic) boundary vertex operators together with a bulk twist operator. This describes the correlator

$$\langle \tau_s \psi \psi \phi \rangle.$$

In canonical picture the vertex operators are

$$\begin{aligned}
V_{\psi_1}^{(-\frac{1}{2})}(z_1) &= \lambda_1 \; e^{-\frac{1}{2}\phi(z_1)} S^{\pm}(z_1) e^{iq_1 \cdot H}(z_1) e^{ik_1 \cdot X}(z_1) \\
V_{\psi_2}^{(-\frac{1}{2})}(z_2) &= \lambda_2 \; e^{-\frac{1}{2}\phi(z_2)} S^{\mp}(z_2) e^{iq_2 \cdot H}(z_2) e^{ik_2 \cdot X}(z_2) \\
V_{\phi}^{(-1)}(z_3) &= \lambda_3 e^{-\phi(z_3)} e^{iq_3 \cdot H}(z_3) e^{ik_3 \cdot X}(z_3) \\
V_{\mathrm{tw}}^{(-1,-1)}(w, \bar{w}) &= \gamma_\theta \; e^{-\phi(w)} e^{-\tilde{\phi}(\bar{w})} \prod_{j=1}^{3} \sigma_{\theta_i}(w, \bar{w}) \\
& \quad \times e^{iq_4 \cdot H(w)} e^{iq_5 \cdot \tilde{H}(\bar{w})} \; e^{ik_4 \cdot X(w)} e^{ik_4 \cdot X(\bar{w})}.
\end{aligned}$$

Picture change, use $SL(2, \mathbb{R})$ symmetry to fix $(w, \bar{w}) \to (i, -i)$ and one of z_1, z_2, z_3 to ∞, and evaluate the 4-pt scattering amplitude.

Quantum Correlator

There are three components, expressed via the useful integral (cf Garousi+Myers 9603194)

$$
\begin{aligned}
I(a,b,c,d,e,f) &= (2i)^f \int_{-\infty}^{\infty} dx_1 \int_{x_1}^{\infty} dx_2 (x_1 - i)^a (x_1 + i)^b (x_2 - i)^c (x_2 + i)^d (x_2 - x_1)^e \\
&= -(2i)^{3+a+b+c+d+e+f}\Gamma(-2-a-b-c-d-e) \times \\
&\Big[(-i)^{2(a+c)} \frac{\sin \pi(b+d+e)\Gamma(1+e)\Gamma(2+b+d+e)\Gamma(-1-d-e)}{\Gamma(-d)\Gamma(-a-c)} \times \\
&{}_3F_2(-c, 1+e, 2+b+d+e; 2+d+e, -a-c; 1) \\
&+(-i)^{2(a+c+d+e)} \frac{\sin(\pi b)\Gamma(-1-c-d-e)\Gamma(1+b)\Gamma(1+d+e)}{\Gamma(-c)\Gamma(-1-a-c-d-e)} \times \\
&{}_3F_2(-d, -1-c-d-e, 1+b; -d-e, -1-a-c-d-e; 1) \Big]
\end{aligned}
$$

Quantum Correlator : First component

Move $\mathcal{V}_\phi(z_3)$ to $z_3 = \infty$, integrate over z_1, z_2:

$$
\begin{aligned}
\mathcal{A}_{z_3\to\infty} \quad : \quad & (u+t) \times I\Big(-1+\theta_1+u/2, -\theta_1+u/2, \theta_2+t/2, 1-\theta_2+t/2, -1+s, -1\Big) \\
& -\frac{t}{2} \times I\Big(-1+\theta_1+u/2, -\theta_1+u/2, -1+\theta_2+t/2, 1-\theta_2+t/2, s, -1\Big) \\
& -\frac{t}{2} \times I\Big(-1+\theta_1+u/2, -\theta_1+u/2, \theta_2+t/2, -\theta_2+t/2, s, -1\Big) \\
& -I\Big(-1+\theta_1+u/2, -\theta_1+u/2, \theta_2+t/2, 1-\theta_2+t/2, -1+s, -1\Big) \\
& +2i\theta_3 \times I\Big(-1+\theta_1+u/2, -\theta_1+u/2, \theta_2+t/2, -\theta_2+t/2, -1+s, -1\Big).
\end{aligned}
$$

Quantum Correlator : Second component

Move $\mathcal{V}_{\psi_2}(z_2)$ to $z_2 = \infty$, integrate over z_1, z_3:

$$\begin{aligned}
\mathcal{A}_{z_2 \to \infty} \quad : \quad & \frac{t}{2} \times I\Big(\theta_3 + s/2, -\theta_3 + s/2, \theta_1 + u/2, -\theta_1 + u/2, -1 + t, -1\Big) \\
& + \frac{t}{2} \times I\Big(\theta_3 + s/2, -\theta_3 + s/2, -1 + \theta_1 + u/2, 1 - \theta_1 + u/2, -1 + t, -1\Big) \\
& + u \times I\Big(\theta_3 + s/2, -\theta_3 + s/2, -1 + \theta_1 + u/2, -\theta_1 + u/2, t, -1\Big) \\
& - I\Big(\theta_3 + s/2, -\theta_3 + s/2, -1 + \theta_1 + u/2, -\theta_1 + u/2, t, -1\Big) \\
& - 2i\theta_3 I\Big(-1 + \theta_3 + s/2, -\theta_3 + s/2, -1 + \theta_1 + u/2, -\theta_1 + u/2, t, -1\Big).
\end{aligned}$$

Quantum Correlator : Third component

Move $\mathcal{V}_{\psi_1}(z_1)$ to $z_1 = \infty$, integrate over z_2, z_3:

$$\begin{aligned}
\mathcal{A}_{z_1\to\infty} \quad : \quad & \frac{t}{2} \times I\Big(-1 + \theta_2 + t/2, 1 - \theta_2 + t/2, \theta_3 + s/2, -\theta_3 + s/2, -1 + u, -1\Big) \\
& + \frac{t}{2} \times I\Big(\theta_2 + t/2, -\theta_2 + t/2, \theta_3 + s/2, -\theta_3 + s/2, -1 + u, -1\Big) \\
& - u \times I(\theta_2 + t/2, 1 - \theta_2 + t/2, \theta_3 + s/2, -\theta_3 + s/2, -2 + u, -1) \\
& + I\Big(\theta_2 + t/2, 1 - \theta_2 + t/2, \theta_3 + s/2, -\theta_3 + s/2, -2 + u, -1\Big) \\
& - 2i\theta_3 I\Big(\theta_2 + t/2, -\theta_2 + t/2, -1 + \theta_3 + s/2, -\theta_3 + s/2, -1 + u, -1\Big).
\end{aligned}$$

Quantum Correlator : Result

The full amplitude comes from combining these three components with appropriate phase factors:

$$\mathcal{A}_{full} = \mathcal{A}_{z_3 \to \infty} - e^{-2\pi i \theta_1} \mathcal{A}_{z_1 \to \infty} - e^{-2\pi i (\theta_1 + \theta_2)} \mathcal{A}_{z_2 \to \infty}.$$

This can be expanded to give poles and finite parts.

For more details and explanations of the subtleties see Lukas Witkowski's talk (3.50pm Tuesday).

Classical Correlator

For a single twist field, classical contribution to

$$\langle \tau_s \psi \psi \phi \rangle$$

vanishes for non-trivial solutions.

Monodromy conditions require

$$\begin{aligned} \partial X(z) &= \alpha(z)(z - z_1)^{-1+\theta}(z - \bar{z}_1)^{-\theta}, \\ \partial \bar{X}(z) &= -\alpha^*(z)(z - z_1)^{-\theta}(z - \bar{z}_1)^{-1+\theta}. \end{aligned}$$

and the action

$$S = \frac{1}{2\pi\alpha'} \int d^2 z \, (\partial X \cdot \bar{\partial} \bar{X} + \partial \bar{X} \cdot \bar{\partial} X)$$

is not normalisable.

Classical Correlator

Now consider classical contributions to the correlator

$$\langle \tau_s^{(1)}(w_1)\tau_s^{(2)}(w_2)...\tau_s^{(n)}(w_n)\psi(z_1)\psi(z_2)\phi(z_3)\rangle.$$

These come from worldsheets stretching between the D3 brane at z_i and the location of the blow-up singularities at w_i.

The classical worldsheet obeys the equation of motion

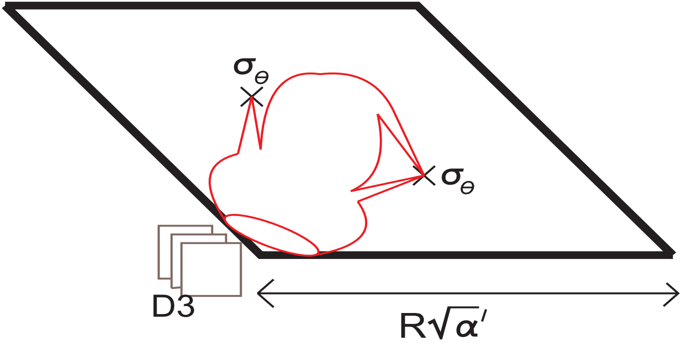
$$\partial\bar{\partial}X = 0,$$

and gives contribution weighted by $e^{-S_{cl}}$ where

$$S_{cl} = \frac{1}{2\pi\alpha'}\int \left(\partial X \cdot \bar{\partial}\bar{X} + \partial\bar{X} \cdot \bar{\partial}X\right)$$

Multi-Twist Correlator

For generic insertions the worldsheet has to stretch



leading to finite-area worldsheets and classical amplitudes suppressed by $e^{-\lambda \frac{R^2}{\alpha'}}$.

Entirely negligible at large volume!

Multi-Twist Correlator

For generic insertions w_i correlator

$$\langle \tau_s^{(1)}(w_1)\tau_s^{(2)}(w_2)...\tau_s^{(n)}(w_n)\psi(z_1)\psi(z_2)\phi(z_3)\rangle.$$

is exponentially suppressed as $e^{-R^2/\alpha'}$ and negligible at LARGE volume.

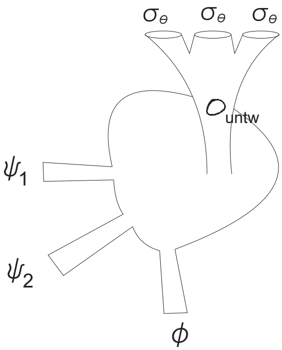
Exponential suppressions originates in local behaviour near twist field

$$X^i(\tau,\sigma) = e^{2\pi i\theta}X^i(\tau,\sigma+2\pi).$$

forcing worldsheet to stretch.

Non-zero correlators can be obtained by limits in which twist fields collide $w_i \rightarrow w_j$ and factorise onto states in the untwisted sector.

Multi-Twist Correlator

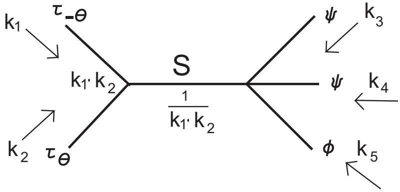


Twist operators factorise onto either massless or massive modes in the untwisted sector.

Multi-Twist Correlator

Factorisation onto massless modes such as the dilaton can lead to a non-zero correlator at zero momentum.

However by construction this factorisation is described by lower-point interactions in the effective field theory.



Interaction is captured by a Lagrangian e.g.

$$\mathcal{L} = S\partial_\mu \tau_\theta \partial^\mu \bar{\tau}_\theta + S\psi\psi\phi.$$

This provides no evidence for the Lagrangian term

$$\mathcal{L} = \tau_s^{(1)} \tau_s^{(2)} \ldots \tau_s^{(n)} \psi\psi\phi.$$

Multi-Twist Correlator

Factorisation onto massive modes such as bulk KK modes requires a different argument.

For an amplitude

$$\langle \sigma^1_{(-1,-1)}(w_1)\sigma^2_{(-1,-1)}(w_2)\ldots\sigma^n_{(-1,-1)}(w_n)\psi_{-1/2}(z_1)\psi_{-1/2}(z_2)\phi_{-1}(z_3)\rangle.$$

picture change all twist operators to $(0,0)$ except first two,

$$\sigma_{(-1,-1)}(w_1) \to \sigma_{(0,-1)}(w_1), \qquad \sigma_{(-1,-1)}(w_2) \to \sigma_{(0,-1)}(w_2).$$

Significance lies in stress tensor OPEs, $e^{\phi} T_F(z) = e^{\phi}(\partial X \cdot \bar{\psi} + \partial \bar{X} \cdot \psi)$

$$\begin{array}{lcl} e^{\phi}\partial\bar{X}\cdot\psi(z) \quad e^{-\phi}\sigma_+ s_+ \tilde{s}_+(w) & \sim & (z-w), \\ e^{\phi}\partial X\cdot\bar{\psi}(z) \quad e^{-\phi}\sigma_+ s_+ \tilde{s}_+(w) & \sim & 1. \end{array}$$

Purely internal picture changing leads to extra 4 uncancellable units of H-charge.

Picture changing must proceed (twice) in external directions.

Multi-Twist Correlator

Factorisation onto massive modes:

Correlator before picture changing:

$$\langle \sigma^1_{(-1,-1)}(w_1)\sigma^2_{(-1,-1)}(w_2)\ldots\sigma^n_{(-1,-1)}(w_n)\psi_{-1/2}(z_1)\psi_{-1/2}(z_2)\phi_{-1}(z_3)\rangle.$$

External picture changing requires the contraction

$$e^{\phi}\partial X\cdot\bar{\psi} \quad e^{-\phi}e^{ik\cdot X}\sim(k\cdot\bar{\psi})e^{ik\cdot x}$$

and involves powers of external momenta.

This forces a momentum prefactor of $k_i \cdot k_j$.

Amplitude vanishes at zero momentum unless we obtain a propagator $\frac{1}{k_i\cdot k_j}$ from factorisation: but factorising onto massive modes cannot give a massless propagator.

$$\longrightarrow \text{No zero momentum term } \int d^4x\,\tau_s^{(1)}\tau_s^{(2)}\ldots\tau_s^{(n)}\psi\psi\phi \text{ in } \mathcal{L}$$

Conclusions

The absence of the zero momentum Lagrangian term

$$\int d^4x \, \tau_s^{(1)} \tau_s^{(2)} \ldots \tau_s^{(n)} \psi\psi\phi$$

shows that physical Yukawas have no dependence on distant blow-up moduli.

By vevving the moduli τ_s we can extend this argument to the smooth resolved space.

The orbifold background also shows that the physical Yukawas do not depend on the bulk Kähler moduli.

At CFT disk level, this establishes sort-of sequestering for these moduli.

Conclusions

Have computed open-closed correlators of form

$$\langle \tau_s^{(1)}(w_1)\tau_s^{(2)}(w_2)...\tau_s^{(n)}(w_n)\psi(z_1)\psi(z_2)\phi(z_3)\rangle.$$

For $n = 1$ we have computed full quantum and classical correlator, for $n > 1$ we establish sort-of sequestering of distant moduli.

Twist moduli at same singularity as D3 branes are not sequestered, twist moduli from distant singularities are.

Next year in Cambridge...

String Pheno 2012: June 25th - 29th Cambridge UK