REMARKS ON BRANES, FLUXES, AND SOFT SUSY BREAKING∗

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We review recent work identifying soft SUSY-breaking terms in local type II string models with branes and magnetic fluxes. We then make a new observation about the configuration space of D-branes in Calabi-Yau backgrounds, and identify vevs for nonperturbative charged hypermultiplets in Calabi-Yau backgrounds with \( N = 2 \) Fayet-Iliopoulos terms.

1. Introduction

A wide class of phenomenologically attractive string theory backgrounds with low-energy \( \mathcal{N} = 1 \) SUSY are described by combinations of D-branes, orientifold planes, and magnetic fluxes. Nontrivial gauge dynamics is typically localized in regions of the compactification manifold, and a fairly generic scenario for SUSY breaking will have supersymmetry broken in one region of the manifold, with the standard model dynamics localized somewhere else. Supersymmetry breaking will be communicated via 10d supergravity effects at tree level, and via radiative corrections as in anomaly or gaugino mediation. However, the detailed appearance and origins of such terms in the low energy effective action of specific models is understood only in a few very specific examples.

In these proceedings we review and slightly extend recent work\(^1\) studying the appearance of tree-level soft SUSY-breaking terms for local models of Calabi-Yau threefold backgrounds with D-branes. The closed string modes live in multiplets of \( \mathcal{N} = 2, \ d = 4 \) SUSY. D-branes and magnetic

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fluxes break SUSY to $\mathcal{N} = 1$ or $\mathcal{N} = 0$. The D-brane modes, and some closed string modes controlling the local geometry, have finite 4d kinetic terms even when the CY is noncompact. The other closed string modes appear as “spurions”, as their dynamics decouples from the low energy physics of the local model. The auxiliary fields for these modes appear as soft SUSY-breaking couplings in the local model. This provides a set of building blocks for more complete models, and in a given model should allow one to address questions such as whether the squark masses are aligned with the quark masses. Therefore, in Section 2 we review the identification of auxiliary fields for light closed string modes with magnetic fluxes, and review how these appear as soft SUSY-breaking couplings on D-branes. This section is based on the talk given by the first author at the Quantum Theory and Symmetries 3 conference at the University of Cincinnati.

Magnetic fluxes and D-branes are crucial aspects of type I, type II, and F-theory models. The interplay between these two aspects of string theory makes apparent some features of the string theory models which are highly nontrivial from the point of view of the low energy field theory. In section 3 we discuss the impact of fluxes on the space of D-branes in Calabi-Yau compactifications. In section 4 we discuss the degrees of freedom responsible for tuning $\mathcal{N} = 2$ Fayet-Iliopoulos terms in type II models on CY threefolds.

Because of the limited space allowed for these proceedings, we will be minimalist about referencing. A more complete bibliography appears in 1. We apologize to those who are not referenced here.

2. Closed string modes and soft SUSY breaking

2.1. Auxiliary fields for $\mathcal{N} = 2$ vector multiplets

We begin with type II string theory on a local (e.g. noncompact) Calabi-Yau threefold $X$ times $\mathcal{R}^4$, together with D-branes filling $\mathcal{R}^4 \times C$, $C \subset X$. Although the D-branes preserve at most $\mathcal{N} = 1$ supersymmetry at most, the closed string modes lie in $\mathcal{N} = 2$ supermultiplets. The $\mathcal{N} = 2$ properties of the latter constrain their couplings to the D-branes 3. Therefore, it is important to understand the underlying $\mathcal{N} = 2$ SUSY structure. Vevs for auxiliary fields in the closed string supermultiplets can break SUSY to $\mathcal{N} = 1$ or $\mathcal{N} = 0$ via computable operators at tree level.

Closed string vector multiplets arise from complex structure deformations of $X$. We can write them as chiral superfields in terms of the $\mathcal{N} = 2$ superspace variables $\theta, \bar{\theta}$ which are a doublet of Weyl spinors under the $SU(2)_R$ symmetry of $\mathcal{N} = 2$ theories. Translations in these superspace
directions are generated by spacetime supercharges built from left- and right-moving worldsheet sectors, respectively.

A vectormultiplet can be described by a superfield $V$ which solves the chiral constraints
\[ \bar{\nabla}_{\dot{\alpha}} V = 0 \]
\[ \hat{\nabla}_{\dot{\alpha}} V = 0 . \]  

The superspace expansion for such a field is:
\[ V = w^a + \theta^\alpha \zeta^a_{\dot{\alpha}} + \theta^2 D^a_{++} + \theta^3 \hat{\theta}^\beta \left( \epsilon_{\alpha\beta} D^a_{+-} + F^a_{\alpha\beta} \right) \]
\[ + \hat{\theta}^2 D^a_{--} + \theta^3 \hat{\theta}^\alpha \chi^a_\alpha + \hat{\theta}^3 \theta^\beta \chi^a_\beta \]
\[ + \theta^2 \hat{\theta}^2 C^a . \]  

One may impose the additional constraints: \(^1, 4\)
\[ (\epsilon_{ij} \nabla^i \sigma_{\mu
u} \nabla^j)(\epsilon_{kl} \nabla^k \sigma^{\mu\nu} \nabla^l)V = -96 \partial^2 \bar{V} \]  

which render $C, \chi$ as dependent variables; impose the constraint that $\sigma^a_\alpha^\beta F^{a\alpha\beta}$ be an anti-self-dual tensor; and impose the “reality constraints” $\Box D^a_{++} = \Box D^a_{--}, \Box D^a_{+-}, \Box D^a_{-+}$ real.

In type IIB string theory, we can identify the bosonic degrees of freedom as follows. If we choose “CFT coordinates” on the moduli space \(^1\), the scalar component $w^a$ can be associated to the perturbation
\[ \delta^m (ds)^2 = \delta g^m_{ij} d\bar{z}^i d\bar{z}^j . \]  

The label $m$ denotes a direction in the complex structure moduli space. Factoring out reparameterizations, each such deformation can be associated to an elements of $H^{(2,1)}(X)$:
\[ \omega^m_{ijk} = \delta g^m_{ij} \delta \bar{g}^{\bar{j}k} \Omega_{ijk} \]  

where $\Omega$ is the holomorphic $(3,0)$ form on $X$ and $g$ is the metric. Choose a basis $\omega^a$ of harmonic representatives of $H^{(2,1)}(X)$. Auxiliary fields correspond to deformations of the NS-NS 3-form $H = \sum_m h^m \omega^m + \text{h.c.}$; of the RR 3-form $F = \sum_m f^m \omega^m + \text{h.c.}$; and of $T = i(\partial - \bar{\partial})J = \sum_m \tau^m \omega^m + \text{h.c.}$, where $J = g_{ij} d\bar{z}^i d\bar{z}^j$. In terms of components $h, f, \tau$:
\[ D^m_{++} = (\tau^m + h^m) \]
\[ D^m_{+-} = g_s (f^m - C^{(0)} h^m) \]
\[ D^m_{--} = (\tau^m - h^m) , \]  

where $g_s$ is the string coupling constant.
where $C^{(0)}$ is the type IIB RR axion.

The results above can be proven using RNS worldsheet techniques. Using these, one may also find the auxiliary fields for the “special geometry” coordinates on complex structure moduli space. That is, choose a symplectic basis $A^a, B_a$ of $H_3(X)$, such that $A^a \cup B_b = \delta^a_b, A^a \cup A^b = B_a \cup B_b = 0$. A good set of coordinates on moduli space is $t^a = \int A^a \Omega$. The “dual periods” $F_a = \int B_a \Omega$ can be written as functions of $t$. The auxiliary fields corresponding to $t^a$ can be written as:

$$
\begin{align*}
D^{a+} &= \int A^a \left( \tilde{T} + \tilde{H} \right) \\
D_{a-} &= g_s \int A^a \left( \tilde{F} - C^{(0)} \tilde{H} \right) \\
D_{a-} &= \int A^a \left( \tilde{T} - \tilde{H} \right),
\end{align*}
$$

(7)

where the tildes denote the projection of the forms into $H^{(2,1)}(X)$. The auxiliary fields for $F_a$ are as above, only with $A^a$ replaced by $B_a$. It is possible to combine these statements into a ‘supermultiplet of three-forms’ which incorporates all of the complex structure multiplets, of the form

$$V = \Omega + \theta^i \theta^j D_{ij} + ...$$

(8)

where $(i, j)$ run over $SU(2)$ doublet indices $\pm$.

A similar story holds for hypermultiplets in type IIA compactifications. 1 The identification of auxiliary fields for vector multiplets in IIA and hypermultiplets in IIB is not yet completely understood.

### 2.2. Soft SUSY breaking

Vevs for auxiliary fields $D_{ij}$ break supersymmetry to $\mathcal{N} = 1$ or $\mathcal{N} = 0$. For example, let $D^{m-} \neq 0$. The SUSY transformations related to $\theta$ are broken, as

$$\delta \tilde{\zeta}_a = \tilde{\epsilon}_a D_{-}.$$  

(9)

If in addition $D_{+-} = D_{++} = 0$, an $\mathcal{N} = 1$ SUSY is still unbroken.

When $D_{ij}$ is related by SUSY to the nonpropagating complex structure deformations of $X$, one may fix its value by hand. In this case, SUSY is explicitly broken by couplings of these nondynamical fields to the propagating modes. One can show explicitly $^{5,1}$ that nontrivial vevs for $G = F - \tau H$, where $\tau = C^{(0)} + i/g_s$, breaks the supersymmetry generating translations along $\theta - i \bar{\theta}$, and leads to a superpotential for complex structure moduli $^{6,2}$. 
We can also use fluxes to introduce soft SUSY breaking terms in $\mathcal{N} = 1$ models with D-branes placed in CY backgrounds. For example, let us study a D5-brane in type IIB wrapping a rational curve $C$ inside $X$, which preserves $\mathcal{N} = 1$ SUSY. Holomorphic deformations of $C$ correspond to open string chiral multiplets, with superfield description $\Phi = \phi + \theta \psi + \theta^2 F_\phi$. To all orders in string perturbation theory, the superpotential

$$W = W(t^a, \Phi^i) = \sum_n g_n(t^a) \text{tr} \Phi^n$$

for these modes depends only on the complex structure moduli of $X$, and not on the Kähler class. If $D_{ij} \neq 0$ is chosen so that the $\mathcal{N} = 1$ SUSY preserved by the D5-brane is broken, one induces explicit, computable SUSY-violating operators. For example, expand $t^a$ in the superspace direction for which the D-brane preserves translation invariance:

$$V_a = t^a + \tilde{\theta}^2 F_a + \ldots$$

where $F$ is the corresponding auxiliary field. The couplings $G_k$ should be written as superfields, so that:

$$g_k \rightarrow g_k(t^a) + \tilde{\theta}^2 F^a \partial_a g_k \equiv g_k + \tilde{\theta}^2 \Delta_k ,$$

leading to soft SUSY-violating terms of the form

$$\int d^2 \theta W + h.c. = \Delta_2 \text{tr} \phi^2 + \Delta_3 \text{tr} \phi^3 + h.c. + \ldots$$

such terms are induced in the presence of RR flux through cycles whose periods appear in the functions $g_k(t^a)$.  

3. Connecting closed-string vacua by paths in open-string field space

In this section, we will show that by moving in open-string configuration space, it is possible to connect vacua with different values of closed-string three-form flux. This amplifies and applies some remarks made in Ref. 7.

Consider type IIB on a CY $X$, with a D5-brane wrapped on a holomorphic curve $C \subset X$ that is a member of a family $\mathcal{M}$ of holomorphic curves such that $\pi_1(\mathcal{M})$ is nontrivial. Examples arise when $C$ is an exceptional curve in the resolution of an $A_1$ singularity over a Riemann surface $M = S_g$ of genus $g > 0$.

The moduli space of D5-branes is lifted by deforming the complex structure of $X$ in such a way that the family $\mathcal{M}$ becomes obstructed; such deformations are in correspondence with sections $dW_0$ of the canonical bundle $T^*S_g$. After a generic such deformation, the moduli space is reduced to a
collection of isolated points where \( 0 = dW_0 \). We will be discussing motion in off-shell configuration space where \( W'_0 \) is not necessarily zero. We may choose \( W_0 \) to be proportional to some small control parameter \( \epsilon \). The potential hills between vacua are then parametrically small compared to the string scale, and the field space for low-energy excitations of the D5-brane is still well-described by \( \mathcal{M} \).

This correspondence between one-forms on the moduli space and complex-structure moduli of the CY implies a map from one-cycles of the moduli space \( \mathcal{M} \) to three-cycles of the CY. \(^{12,11}\) A path \( \gamma \) maps to the three-cycle \( \pi^{-1}\gamma \) obtained by fibering the exceptional curve \( C_x \) over each point in \( \gamma \). Moving the D-brane around a loop \( \gamma \) in the moduli space \( \mathcal{M} \) generates a quantum of RR flux though the cycle \( \pi^{-1}\gamma \). \(^7\) This follows from the fact that the D5-brane is magnetically charged under the three-form flux. \(^8\) We can see this fact further manifest itself in the superpotential. \(^9\)

The effective superpotential governing the open-string moduli and complex structure moduli is: \(^{8,6,9}\)

\[
W = W_{GVW} + W_{obstruction} = \int_X \Omega \wedge G + \int_{\Xi} \Omega 
\]  

where \( G \) is the three-form associated to a linear combination of \( D_{ij} \) preserving the same SUSY as the D5-brane, and includes the RR flux. \(^1\) For deformations of a D5-brane from a rational curve \( C_0 \) to a curve \( C \), the three-chain \( \Xi = C - C_0 \). There are two ambiguities in defining this obstruction contribution to the superpotential. \(^8\)

1. \( C_0 \) is a base point on the closed string moduli space; changing \( C_0 \) changes \( \Xi \) and so changes \( W \) by an additive constant.
2. Since \( H_3(X) \) is nontrivial, a 3-chain \( \Xi \) such that \( \partial \Xi = C - C_0 \) is only determined up to the addition of an element of \( H_3(X, \mathbb{Z}) \). This also additively changes the superpotential. We will give a concrete example below.

Moving the brane in a loop \( \gamma \) in \( \mathcal{M} \) shifts the chain \( \Xi \) by \( \Xi \mapsto \Xi + \pi^{-1}\gamma \). This in turn shifts the superpotential by \( \int_{\pi^{-1}\gamma} \Omega \). By Poincaré duality, this can be identified with

\[
\delta W = \int_X \Omega \wedge [\pi^{-1}\gamma]. 
\]

\(^a\)An illustrative analogy arises in Maxwell theory on \( R^3 \times S^1 \): start with a magnetic monopole-antimonopole pair, and move the magnetic monopole around the circle. This causes the magnetic flux through the transverse plane to jump.
But this can be absorbed in $W_{GV}$, if the RR flux shifts by $[\pi^{-1}\gamma]$. This possibility of interchanging contributions between the two terms in $W$ is made clearest by writing

$$W = \int_{G+\Xi} \Omega. \quad (15)$$

where $G$ is the 3-cycle Poincaré dual to the flux.

Let us consider an explicit example of a patch of this model. Consider the hypersurface in $\mathbb{C}^4$ given by

$$y^2 + u^2 + v^2 = W'_0(x)^2 + f(x). \quad (16)$$

In this example, we take $C$ to be an $S^2$ which can be resolved out of any double roots of the RHS of (16) at a point $x$, and $\Xi$ is this $S^2$ times a curve in the $x$-plane ending at $x$. There is no reason for $\Xi$ to be special Lagrangian, and it is not. The superpotential is $W = \int_\Xi \Omega$. All of the information about the threefold and its complex structure, including this integral, can be represented in terms of information on the Riemann surface $\Sigma$ at $u = v = 0$, defined by $y^2 = (W'_0)^2 + f$. $\Sigma$ is a double cover of the $x$-plane, each of whose fibers represent a two-sphere homologous to $C$. The superpotential integral can be represented as

$$W(x) = \int_{x_0}^{x} y(\tilde{x})d\tilde{x} \quad (17)$$

where $C$ is the $S^2$ over the point $x$, and $x_0$ specifies the base-point curve.

Fig. 1: Each point in this picture represents a hemisphere of the exceptional $P^1$ in the ALE singularity $y^2 + u^2 + v^2 = 0$.

It is important to distinguish $\Sigma$ from $S_g$. The geometry of the $x$-plane embeds into $S_g$ as shown in Fig. 1. In this example, the moduli space of the
curve $C$ when $W_0' = 0, f = 0$ is $S_g$; for generic $W_0$ at $f = 0$, holomorphic curves appear only at the critical points of $W_0$.

The shift in the superpotential we have described is effected by changing the flux through the cycle dual to the 3-cycle the D5-brane sweeps out in moving through a loop in $H_1(M)$. We can see further that this is consistent with rules for tadpole cancellation when one turns on $N$ units of NS-NS three-form flux $H$ through the three-cycle $Ξ$ that the D5-brane is sweeping out.

In this case, if the D5-brane sweeps out a cycle, we have stated that it induces a jump $δF = [π^{-1}γ]$ in the RR flux. Since $\int H \wedge F = M$ induces an RR 4-form tadpole that must be cancelled by adding $M$ three-branes. But this tadpole is precisely cancelled by D3-brane charge on the D5-brane which sweeps out the cycle $Ξ$ with H-flux. H-flux through $Ξ$ means that there is a gradient for the B-field through the sphere the D5-brane is wrapping, with respect to the direction on $M$ it is moving. H-flux quantization means that the B-field will shift by $2πN$ upon traversing the loop $γ$, in units where $B = B + 2π$ when there is no brane. Because of this, $B$ induces $N$ units of D3-charge via the worldvolume Chern-Simons coupling $\int_{D5} B \wedge C(4)$. This phenomenon is essentially identical to the phenomenon described in $^{13}$, domain walls in $R^4$ achieved by wrapping D5-branes around cycles $Ξ$ with NS-flux $\int_Ξ H = N$ interpolate between vacua with D3-brane number differing by $N$.

This result implies that that in going around what was apparently a loop in this open-string moduli space $M$, the string theory does not come back to itself. Rather, the closed-string background is changed. We conclude that the open-string moduli space in fact has no $π_1$. For example, stable cosmic string solutions corresponding to the putative loop do not exist. One indication is that if one turns on the obstruction superpotential, any loop in $R^4$ for which the D-brane position loops around a cycle in $S_g$ will cross a domain wall. $^{14}$ This domain wall is a D5-brane wrapping the three-cycle $Ξ$. The jump in RR flux induced by the motion in moduli space as one circles the cosmic string is then cancelled by the jump in flux induced by the domain wall, so that the 4d solution is single-valued. In fact this entire domain-wall-ending-on-cosmic-string is one boundary-less D5-brane.
Fig. 2: The vertical direction $v$ is the would-be cycle in the configuration space; hashed lines are identified. The blue lines represent the D5-brane on $C$ whose $v$ position depends on the argument of $x + iy$. The red lines represent the domain wall D5-brane wrapping $\Xi$. Note that their boundaries cancel. This figure makes the instability of the configuration clear: along the half-line where the brane crosses itself (indicated by the wavy line), it can annihilate; it can subsequently slip off the $v$-circle.

Said another way, the moduli space or low-energy field space is in fact a multiple cover – with infinitely many sheets – of the Riemann surface $S_g$, and the RR flux labels the sheets. Note that this generically does involve going off-shell, since in the presence of the obstruction superpotential there is not in fact a moduli space.

A similar discussion implies that it is possible to interpolate between values of NSNS flux quanta by moving wrapped NS5-branes.

4. $\mathcal{N} = 2$ Fayet-Iliopoulos terms in type II models

Taylor and Vafa \(^2\) showed that in local models of type IIB Calabi-Yau compactifications, the superpotential \(^6\) for complex structure moduli can arise from electric and magnetic Fayet-Iliopoulos (FI) terms \(^{15}\) which spontaneously break the global $\mathcal{N} = 2$ SUSY to $\mathcal{N} = 1$. This is consistent with the identification of magnetic flux \(^{1,5}\) with auxiliary fields, as the auxiliary fields will be equated to the FI terms on-shell. \(^{15}\)

One might ask whether the FI terms should be identified as separate degrees of freedom, equated to the magnetic flux via on-shell equations of motion. This appears to be the case. Study a deformed conifold in type IIB
with vanishing 3-cycle $A$ and dual cycle $B$. D3-branes wrapping $A$ are light hypermultiplets charged with respect to the vector multiplets associated with the period $t_A$. The hypermultiplet can be written in terms of two $\mathcal{N} = 1$ chiral multiplets $Q, \tilde{Q}$ with opposite $U(1)$ charge. If we write the $\mathcal{N} = 2$ supermultiplet $V = t_A + \cdots$ in terms of an $\mathcal{N} = 1$ chiral multiplet $A$ and a $\mathcal{N} = 1$ vector multiplet, the coupling of $Q, \tilde{Q}$ to $A$ includes the following superpotential term:

$$W = \int d^2 \theta A Q \tilde{Q}$$

If the scalar component of $< Q\tilde{Q} >$ gets a vev, this will appear as an electric FI term, or a magnetic FI term with respect to the vector multiplet associated to $t_B$.

It would be nice to show microscopically that such a vev induces magnetic 3-form flux, and that the potential for $< Q\tilde{Q} >$ has discrete minima associated to different values of NS-NS and RR flux through $B$. We can, however, note that when $\int_A H = -K \gg 1, \int_B F = N \gg 1$, this result is consistent with the conjectured field theory dual. This geometry is described by an $\mathcal{N} = 1$ $SU(NK+N) \times SU(N)$ gauge theory with bifundamentals in $(NK+N, \bar{N})$ and $(NK+N, N)$. At low energies, the gauge invariant degrees of freedom include “meson” and “baryon” degrees of freedom, constructed from the bifundamentals. The mesons are dual to motions of D3-branes in the Klebanov-Strassler geometry. The baryons correspond to D3-branes wrapping the 3-cycles of this geometry. The space of vacua contains branches where either the mesons or baryons have vevs. Domain walls connecting the meson and baryon branches were argued to be dual to D5-branes wrapping $A$. The disappearance of the D3-branes is consistent with the tadpole cancellation arguments reviewed above.

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References

1. A. Lawrence and J. McGreevy, arXiv:hep-th/0401034.
2. T. R. Taylor and C. Vafa, Phys. Lett. B474, 130 (2000) [arXiv:hep-th/9912152].
3. I. Brunner, M. R. Douglas, A. E. Lawrence and C. Romelsberger, J.High Energy Phys. 0008, 015 (2000) [arXiv:hep-th/9906200].
4. M. de Roo, J. W. van Holten, B. de Wit and A. Van Proeyen, Nucl. Phys. 173, 175 (1980).
5. C. Vafa, J. Math. Phys. 42, 2798 (2001) [arXiv:hep-th/0008142].
6. S. Gukov, C. Vafa and E. Witten, Nucl. Phys. B584, 69 (2000) [Erratum-ibid. B608, 477 (2001)] [arXiv:hep-th/9906070].
7. M. Aganagic, A. Klemm and C. Vafa, Z. Naturforsch. A 57, 1 (2002) [arXiv:hep-th/0105045].
8. E. Witten, Nucl. Phys. B507, 658 (1997) [arXiv:hep-th/9706109].
9. W. Lerche, P. Mayr and N. Warner, arXiv:hep-th/0208039; W. Lerche, arXiv:hep-th/0312326.
10. S. Kachru, S. Katz, A. E. Lawrence and J. McGreevy, Phys. Rev. D62, 126005 (2000) [arXiv:hep-th/0006047].
11. S. Katz, D. R. Morrison and M. Ronen Plesser, Nucl. Phys. B477, 105 (1996) [arXiv:hep-th/9601108].
12. C. H. Clemens and P. A. Griffiths, Ann. Math. 95, 281 (1972).
13. S. Kachru, J. Pearson and H. Verlinde, J. High Energy Phys. 0206, 021 (2002) [arXiv:hep-th/0112197].
14. E. J. Copeland, R. C. Myers and J. Polchinski, arXiv:hep-th/0312067.
15. I. Antoniadis, H. Partouche and T. R. Taylor, Phys. Lett. B372, 83 (1996) [arXiv:hep-th/9512066]; I. Antoniadis and T. R. Taylor, Fortsch. Phys. 44, 487 (1996) [arXiv:hep-th/9604062]; H. Partouche and B. Pioline, Nucl. Phys. Proc. Suppl. 56B, 322 (1997) [arXiv:hep-th/9702115].
16. I. R. Klebanov and M. J. Strassler, J. High Energy Phys. 0008, 052 (2000) [arXiv:hep-th/0007191].
17. S. S. Gubser and I. R. Klebanov, Phys. Rev. D58, 125025 (1998) [arXiv:hep-th/9808075].