Spontaneous SUSY Breaking in Various Dimensions

Amit Giveon\textsuperscript{1}, David Kutasov and Oleg Lunin

EFI and Department of Physics, University of Chicago
5640 S. Ellis Av., Chicago, IL 60637, USA

We generalize the ISS model of spontaneous supersymmetry breaking to lower dimensions. We also comment on the dynamics of the corresponding brane systems in string theory, and on possible applications to gauge/gravity duality.

\textsuperscript{1} Permanent address: Racah Institute of Physics, The Hebrew University, Jerusalem 91904, Israel.
1. Introduction

In this note we will discuss a class of gauge theories in $d < 4$ spacetime dimensions that exhibit spontaneous supersymmetry breaking, either in the vacuum or in metastable states. Our main motivation for studying these theories comes from gauge/gravity duality. One way to study supersymmetry breaking in $(d+1)$-dimensional gravity is via a $d$-dimensional gauge theory dual. If the gravitational theory describes a large smooth $(d+1)$-dimensional spacetime, the dual gauge theory is typically strongly coupled, and thus difficult to solve. We will restrict attention to weakly coupled gauge theories, which are easier to analyze, but our results may be useful for studying gauge theories with semiclassical gravity duals.

The models we will study can be introduced by starting with the ISS model of spontaneous supersymmetry breaking in four dimensions [1], and dimensionally reducing it to $d < 4$. This gives a gauge theory with four supercharges, and it is natural to ask whether it breaks supersymmetry, like its four dimensional analog. As in four dimensions, one can choose the parameters of the model such that the gauge interactions are weak and can be neglected. In that regime, the low energy degrees of freedom are chiral superfields, and their dynamics is described by a Wess-Zumino (WZ) model. In four dimensions, this model was studied by ISS; our main purpose is to generalize their results to $d < 4$. We will also discuss the regime of validity of the field theory analysis, and its relation to the dynamics of the brane system in string theory that realizes this WZ model at low energies.

The case $d = 3$ is particularly interesting for gauge/gravity duality, since the corresponding gravitational theory is $3 + 1$ dimensional. In that case one can replace the standard Yang-Mills (YM) kinetic term by a Chern-Simons (CS) one. Many CS matter theories are expected to have $AdS_4$ holographic duals, and thus provide a natural arena for connecting our discussion to four dimensional gravity.

This note is organized as follows. In section 2 we study the WZ models obtained by reducing the (generalized) ISS theory to $d < 4$ dimensions. We show that they have metastable supersymmetry breaking vacua, which can be studied reliably at weak coupling. In section 3 we discuss the gauging of the color group, with either a YM or (for $d = 3$) CS kinetic term for the gauge field. In section 4 we summarize and comment on our results, and the corresponding brane systems in string theory.
2. SUSY breaking in $d < 4$ WZ models

In this section we consider a WZ model in $d$ spacetime dimensions, with chiral superfields $\Phi^i, q^i_a, \tilde{q}^a_i$, where $i, j = 1, \ldots, N_f$, $a = 1, \ldots, N$, and

$$N_f > N. \quad (2.1)$$

As we will see later, when the condition (2.1) is not satisfied, the model does not break supersymmetry, even in metastable states.

The Kähler potential for all the fields is taken to be canonical. The superpotential is given by

$$W = h q \Phi \tilde{q} + h \text{Tr} \left( \frac{1}{2} \epsilon \mu \Phi^2 - \mu^2 \Phi \right). \quad (2.2)$$

The couplings $h, \epsilon, \mu$ are in general complex. Two of them can be taken to be real and positive by redefining $\Phi, \tilde{q}, q$. We will sometimes take all three couplings to have this property, for simplicity. In that case, the interesting dynamics occurs at real $\Phi$. It is easy to repeat the discussion for the general case.

For $\epsilon = 0$, the model (2.2) is obtained by a dimensional reduction of that discussed in [1]. For $\epsilon \neq 0$, it is similarly related to the model studied in [2-5]. As we will see later, a non-zero $\epsilon$ is needed for $d < 4$ to control the dynamics. The model (2.2) has a global symmetry $U(N) \times U(N_f)$. We will eventually be interested in gauging $U(N)$, but for now we treat it as a global symmetry.

The couplings that appear in the superpotential (2.2) have the following mass dimensions:

$$[h] = 2 - \frac{d}{2}, \quad [\mu] = \frac{d}{2} - 1, \quad [\epsilon] = 0. \quad (2.3)$$

For $d < 4$, all terms in the superpotential are relevant. Thus, in general the WZ model (2.2) is expected to be strongly coupled in the infrared. We will see later that this can be avoided when the vacua of interest are located sufficiently far from the origin of field space.

The tree level bosonic potential corresponding to (2.2) is

$$V_0 = |h|^2 \left( |q^i \tilde{q}_j - \mu^2 \delta^i_j + \epsilon \mu \Phi_j^i|^2 + |q \Phi|^2 + |\Phi \tilde{q}|^2 \right). \quad (2.4)$$

Supersymmetric vacua are labeled by an integer $k = 0, \ldots, N$. For given $k$, the expectation values of the chiral superfields take the form (up to global symmetries)

$$\Phi = \begin{pmatrix} 0 & 0 \\ 0 & \frac{\mu}{\epsilon} I_{N_f - k} \end{pmatrix}, \quad q \tilde{q} = \begin{pmatrix} \mu^2 I_k \\ 0 \end{pmatrix}, \quad (2.5)$$
where $I_k$ is a $k \times k$ identity matrix. In a vacuum with given $k$, the $U(N) \times U(N_f)$ symmetry is broken to $U(N - k) \times U(N_f - k) \times U(k)$ by the expectation values of $\Phi, q, \tilde{q}$.

In four dimensions, in addition to the supersymmetric vacua (2.5), the WZ model has metastable vacua in which the classical potential (2.4) is balanced by a one loop, Coleman-Weinberg (CW) contribution [2, 4]. It is natural to ask whether such vacua occur in lower dimensions as well.

Following [2-5], we parameterize the fields as follows:

$$\Phi = \begin{pmatrix} 0 & 0 & 0 \\ 0 & XI_n & 0 \\ 0 & 0 & \frac{\mu}{\epsilon} I_{N_f-k-n} \end{pmatrix}, \quad q\tilde{q} = \begin{pmatrix} \mu^2 I_k & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad (2.6)$$

and examine the effective potential for $X$. The tree level potential (2.4) is

$$V_0(X) = n|h\mu|^2|\mu - \epsilon X|^2. \quad (2.7)$$

The Coleman-Weinberg potential is obtained by integrating out the components of $\Phi, q, \tilde{q}$, whose mass depends on $X$. It is given by

$$V_1(X) = 2nk \sum_{\eta, \sigma = \pm 1} [F_d(m_b) - F_d(m_f)] + 2n(N - k) \sum_{\eta = \pm 1} [F_d(\hat{m}_b) - F_d(\hat{m}_f)] , \quad (2.8)$$

where

$$F_d(m) = \frac{1}{2} \int \frac{d^dp}{(2\pi)^d} \ln(p^2 + m^2) = \frac{|m|}{4}, -\frac{m^2}{8\pi} \ln m^2, -\frac{|m|^3}{12\pi}, \frac{m^4}{64\pi^2} \ln m^2, \quad (2.9)$$

in $d = 1, 2, 3, 4$, respectively.

The first term in (2.8) (proportional to $kn$) is the contribution of the $k \times n$ off-diagonal components of $q, \tilde{q}$ and $\Phi$ that transform in the fundamental of the $U(k)$ subgroup of $U(N) \times U(N_f)$ unbroken by the configuration (2.6). The second is due to the components of $q, \tilde{q}$ that transform in the fundamental of the unbroken $U(N - k)$ subgroup. The masses in (2.8) depend on $X$ as follows [4]:

$$m_b^2 = |h^2| \left( |\mu|^2 + \frac{1}{2} |X|^2 + \frac{|\epsilon \mu|^2}{2} + \frac{\eta G}{2} + \frac{\sigma}{2} \sqrt{(|X|^2 - |\epsilon \mu|^2) + 4|\mu X^* + \epsilon| \mu^2|^2} \right),$$

$$m_f^2 = |h^2| \left( |\mu|^2 + \frac{1}{2} |X|^2 + \frac{|\epsilon \mu|^2}{2} + \frac{\sigma}{2} \sqrt{(|X|^2 - |\epsilon \mu|^2) + 4|\mu X^* + \epsilon| \mu^2|^2} \right),$$

$$\hat{m}_b^2 = |h^2| (|X|^2 + |G + \mu X|), \quad \hat{m}_f^2 = |h^2| (|X|^2) , \quad (2.10)$$
The bosonic masses $m_b$ on the first line of (2.10) in the sector with $\eta = \sigma = -1$ are tachyonic near the origin\footnote{The masses $\hat{m}_b$ on the third line are tachyonic as well for $\eta = -1$ and small $X$; we will return to them later.} and need to be treated more carefully. One can think of the field configuration (2.4) as a deformation of the supersymmetric vacuum, (2.3), with $k \to k + n$. If the effective potential for $q, \tilde{q}, X$ has a non-supersymmetric local minimum of the form (2.6), the global symmetry of the relevant sector of the theory is broken: $U(k + n) \to U(k) \times U(n)$. This leads to $2kn$ Nambu-Goldstone bosons parametrizing the coset $U(k + n)/U(k) \times U(n)$. These are precisely the $2kn$ modes with $\eta = \sigma = -1$ mentioned above. Their masses must receive contributions due to one loop effects that make them zero at the local minimum.

To analyze the one loop effective potential (2.8), consider first the case $k = N$, for which the second term in (2.8) is absent (i.e. the $U(N)$ symmetry is completely broken). Expanding the first term in (2.8) to second order in $X/\mu$ one finds

\[
\begin{align*}
\text{d} = 1 & : & V_1 = \frac{knh}{8} \left[ 4\mu(\sqrt{2} - 2) + (4 + \sqrt{2})\epsilon X + 2(\sqrt{2} - 1)\frac{X^2}{\mu} + \cdots \right], \\
\text{d} = 2 & : & V_1 = -\frac{knh^2}{4\pi} \left[ \epsilon \mu X \ln(h^2 \mu^2) + (2X^2 + 2\mu^2 + \epsilon \mu X) \ln 2 - 3\epsilon \mu X - 2X^2 + \cdots \right], \\
\text{d} = 3 & : & V_1 = \frac{knh^3}{4\pi} \left[ -\frac{4}{3}(\sqrt{2} - 1)\mu^3 + (2 - \sqrt{2})\epsilon \mu^2 X + (3 - 2\sqrt{2})\mu X^2 + \cdots \right], \\
\text{d} = 4 & : & V_1 = \frac{knh^4}{16\pi^2} \left[ \frac{\mu^3(\mu - 2\epsilon X) \ln(h^2 \mu^2)}{4} + (\mu^4 + \epsilon \mu^3 X + 2\mu^2 X^2) \ln 4 - 5\epsilon \mu^3 X - 2\mu^2 X^2 + \cdots \right].
\end{align*}
\]

The contributions from higher orders in the loop expansion and from nonperturbative effects can be neglected when the coupling $h$ at the scale of the masses of the fundamentals $q, \tilde{q}$ is small:

\[
kh^2 \ll (hX)^{4-d}.
\]

The potential $V_0 + V_1$ has a local minimum at\footnote{Here we assume that $\epsilon^2 \ll bk; |a|k \ll 1$; we will justify these approximations shortly.}

\[
X_0 \simeq \frac{\epsilon \mu (1 + ak)}{\epsilon^2 + bk} \approx \frac{\epsilon \mu}{bk},
\]

where $a$ and $b$ are dimension dependent constants.
\[
\begin{array}{|c|c|c|}
\hline
d & a & b \\
\hline
1 & -\frac{4+\sqrt{2}}{16\hbar\mu^4} & \frac{\sqrt{2}-1}{4\hbar\mu^4} \\
2 & \frac{\ln(2h^2\mu^2)-3}{8\pi\mu^2} & \frac{1-\ln 2}{2\pi\mu^2} \\
3 & -\frac{(2-\sqrt{2})\hbar}{8\pi\mu} & \frac{(3-2\sqrt{2})\hbar}{4\pi\mu} \\
4 & \frac{\hbar^2}{32\pi^2} \left(5 + 2 \ln \frac{h^2\mu^2}{2}\right) & \frac{\ln 4-1}{8\pi^2} h^2 \\
\hline
\end{array}
\]

The expression for the position of the local minimum, (2.13), is valid when \( \epsilon \) is in the range

\[bk \gg \epsilon \gg bk \frac{k^\frac{1}{4-\alpha}}{4-\alpha} \mu , \quad (2.14)\]

where the lower bound comes from the condition for validity of the one loop approximation, (2.12), and the upper bound is the condition \( X \ll \mu \), which was used in approximating the Coleman-Weinberg potential by (2.11).

Note that the existence of the range (2.14) implies a lower bound on \( \mu \):

\[\mu^{4-d} \gg k\hbar^{d-2} . \quad (2.15)\]

In this range, \(|a|, bk, \epsilon \ll 1\); this, together with (2.14), provides a justification for the approximations mentioned in footnote 3. To estimate the lifetime of the metastable vacuum, we need to compute the Euclidean action which gives the amplitude of tunneling to the true minimum. Approximating the barrier by a triangular potential and applying the arguments of [3] to the \( d \)-dimensional field theory, we find

\[S \sim V_{\text{met}}^{1-d/2} (X_{\text{susy}})^d \sim (h\mu^2)^{2-d} \left(\frac{\mu}{\epsilon}\right)^d \gg \epsilon^{-d} , \quad (2.16)\]

where in the last inequality we used (2.13). For \( \epsilon \ll 1 \), the non-supersymmetric state at (2.13) is parametrically long-lived, and is at a value of \( X \) which is well separated from the supersymmetric vacuum, located at \( X_{\text{susy}} = \mu/\epsilon \gg \mu \).

We end this section with some comments:

\[\begin{array}{l}
4 \text{ Here and below we restrict to } d < 4 . \\
5 \text{ The case of } d = 3 \text{ is discussed in more detail in [3]} .
\end{array}\]
(a) Validity of the loop expansion in the WZ model requires the one loop potential $V_1$ to be much smaller than the tree level contribution $V_0$. Equations (2.4), (2.11) imply that
\[
\frac{V_1}{V_0} \sim k \left( \frac{h \mu}{(h \mu^2)^2} \right)^d = k \frac{h^{d-2}}{\mu^{1-d}} \ll 1.
\] (2.17)
The last inequality in (2.17) follows from (2.15).

(b) At the local minimum (2.13), the pseudo-moduli $X$ have a mass, which can be read off from (2.11),
\[
m_X^2 = bk(h \mu)^2.
\] (2.18)
Since $bk$ is small in the regime (2.15), the mass (2.18) is well below the masses of the quarks (2.10), which were integrated out in calculating the potentials (2.11). Thus, the above analysis is consistent.

(c) We restricted the discussion to the small field region, where the one-loop potential can be approximated by (2.11), and the effective coupling $h$ is small, (2.12). It is natural to expect that the results can be extended beyond this region. For example, as $\epsilon \to 0$, the supersymmetric vacuum goes to infinity, and the non-supersymmetric one, (2.13), approaches the origin. The lower bound (2.14) is violated, and the low energy dynamics becomes strongly coupled. Although this limit is not under control, we expect the theory with $\epsilon = 0$ to have a stable, non-supersymmetric vacuum at the origin [8].

(d) For $k < N$, some of the components of $q, \tilde{q}$ are tachyonic for small $X$ (see third line of (2.10)). To avoid these tachyons one must go to $X \sim \mu$, or larger. Studying the one-loop corrected potential $V_0 + V_1$ numerically, we find that, as in the four dimensional analysis of [4], there are no local minima of the full potential, which do not decay by condensation of the above tachyonic modes.

3. Gauging $U(N)$

Unlike its four dimensional analog, in $d \leq 3$ the WZ model with superpotential (2.2) is UV complete. However, for applications to gauge/gravity duality it is interesting to consider a generalization of the model in which part or all of the global symmetry group $U(N) \times U(N_f)$ is gauged.

Consider, for example, the case where we gauge $U(N)$ (of course, requiring that the Lagrangian still preserves SUSY). This introduces into the problem a new parameter, the
gauge coupling \(g_{YM}\), whose mass dimension is \([g_{YM}] = 2 - d/2\), the same as that of \(h\), (2.3). The gauge interaction is a relevant perturbation of the WZ model, which grows in the infrared, just like \(h\). The results of the previous section are not modified by its presence when \(g_{YM} \ll h\), so that the effect of gauge theory loops is smaller than that of loops in the WZ model, taken into account in section 2.

The \(U(N)\) gauge field belongs to a supermultiplet that also contains \(4 - d\) real scalar fields in the adjoint representation. In analyzing the non-supersymmetric vacuum structure, one needs to determine the potential of these fields as well as of those considered in section 2. However, since in the WZ model SUSY breaking vacua occur only in the sector with \(k = N\), where the \(U(N)\) symmetry is completely broken, these scalars get a mass of order \(g_{YM}\mu\), and do not destabilize the vacua found in section 2.

In three dimensions, one can also consider a theory in which the kinetic term of the \(U(N)\) gauge field takes the Chern-Simons form. The Chern-Simons level \(k_{cs}\) is quantized, but at large level one can study the gauge dynamics perturbatively in the 't Hooft coupling \(N/k_{cs}\). The analysis of the previous section is valid to leading order in the 't Hooft coupling.

In [8] it was argued that the above CS theory is related by a Seiberg-type duality \(^6\) to \(U(N_c)\) Chern-Simons theory at level \(k_{cs}\), coupled to \(N_f\) fundamentals \(Q_i, \tilde{Q}^i\), with superpotential

\[
W_{el} = m\tilde{Q}Q + \lambda(\tilde{Q}Q)^2.
\]  

(3.1)

The parameters \(\epsilon\) and \(\mu\) in (2.2) are related to \(m\) and \(\lambda\), while the field \(\Phi\) is proportional to the electric meson field \(\tilde{Q}Q\). The rank of the magnetic gauge group \(N\) is given by [8]:

\[
N = N_f + k_{cs} - N_c.
\]  

(3.2)

The condition for supersymmetry breaking (2.7) is \(k_{cs} < N_c\). It is believed [12] that for \(\lambda = 0\) the electric theory spontaneously breaks supersymmetry for \(k_{cs}\) in this range. This is consistent with the fact that the magnetic theory is expected to break SUSY for \(\epsilon = 0\).

4. Discussion

The main conclusion of our analysis is that the appearance of spontaneous supersymmetry breaking metastable vacua, found by ISS [1] in four dimensional gauge theory,

\(^6\) For generalizations of this duality, see e.g. [9,10,11].
persists to lower dimensions as well. For generic values of the parameters, the $d < 4$ dimensional gauge theory with superpotential (2.2) is strongly coupled in the infrared, and the existence of such vacua is difficult to establish. However, there are regions in parameter space where these vacua can be studied reliably at weak coupling. The phase structure we found is similar to that found in four dimensions in [1,2]. For example, in terms of the decomposition (2.5), the branch with $k = N$ contains SUSY breaking vacua, while the “tachyonic branches” with $k < N$ do not.

As mentioned in the introduction, one of the primary motivations for this study is gauge/gravity duality. A useful step in that direction is to realize the systems we studied in string theory. This can be done by using the brane constructions reviewed in [13].

\[ \text{Fig. 1: The brane configuration.} \]

A brane system that realizes the gauge theory described in section 3 is exhibited in figure 1. The $d$ dimensional gauge theory is realized on $D_p$-branes with $p = d$, stretched between $NS5$-branes and other $D$-branes. The parameters of the WZ model (2.2) are related to the brane ones as follows:

\[ h^2 = \frac{2g_s(2\pi)^{p-2}l_s^{-3}}{y_2 - y_1}, \quad \mu^2 = \frac{v_2}{2g_s(2\pi l_s)^{p-1}}, \quad \epsilon^2 = \frac{2\pi l_s^2 \tan^2 \theta}{v_2(y_2 - y_1)}. \] (4.1)

Figure 1 gives rise to a gauge theory with standard kinetic term for the $U(N)$ gauge field. As discussed in section 3, in three dimensions (i.e. for $p = 3$ in figure 1), one can also consider a theory in which the kinetic term has a CS form. In the brane picture this corresponds to replacing the $NS'$-brane by a $(1,k_{cs})$ brane, oriented such that supersymmetry is unbroken [14].
The relation between the low energy dynamics of the branes and the gauge theory in four dimensions has been recently discussed in [5]. The picture in $d < 4$ is very similar. In the region of parameter space where the geometric description of figure 1 is reliable (when all the distances are much larger than $l_s$, and the string coupling $g_s$ is small), the loop effects in the WZ model, (2.11), and the corresponding gauge theory effects, are small. The dominant corrections are due to the gravitational attraction of the $Dp$-branes to the $NS$-brane in figure 1. The balance of this attraction and the classical effects leads to the appearance of metastable vacua, which can be studied as in [15,3].

From the point of view of the low energy theory, these gravitational effects give rise to a non-trivial Kähler potential and higher D-terms for the pseudo-moduli, which together with the non-trivial superpotential stabilize them [3]. The brane picture leads to a phase structure which is essentially the same in all dimensions, since the D-terms that govern it are independent of $d$ when written in terms of geometric variables. Therefore, it is nice that the field theory analysis also gives a phase structure which is similar for different dimensions.

Low energy theories on branes play an important role in gauge/gravity duality. It would be interesting to apply our results in that context. A particularly promising class of theories are three dimensional CS theories which often have supergravity duals. We have seen that for large CS level, $k_{cs} \gg N$, the appearance of supersymmetry breaking vacua is rather generic. It would be interesting to see whether this is also the case in the gravity regime $N \gg k_{cs} \gg 1$.

**Acknowledgements:** We thank G. Torroba for correspondence. This work is supported in part by the BSF – American-Israel Bi-National Science Foundation. AG is supported in part by a center of excellence supported by the Israel Science Foundation (grant number 1468/06), DIP grant H.52, and the Einstein Center at the Hebrew University. DK and OL are supported in part by DOE grant DE-FG02-90ER40560 and the National Science Foundation under Grant 0529954. AG thanks the EFI at the University of Chicago for hospitality during the course of this work.
References

[1] K. A. Intriligator, N. Seiberg and D. Shih, “Dynamical SUSY breaking in meta-stable vacua,” JHEP 0604, 021 (2006) [arXiv:hep-th/0602239].

[2] A. Giveon and D. Kutasov, “Stable and Metastable Vacua in SQCD,” Nucl. Phys. B 796, 25 (2008) [arXiv:0710.0894 [hep-th]].

[3] A. Giveon and D. Kutasov, “Stable and Metastable Vacua in Brane Constructions of SQCD,” JHEP 0802, 038 (2008) [arXiv:0710.1833 [hep-th]].

[4] R. Essig, J. F. Fortin, K. Sinha, G. Torroba and M. J. Strassler, “Metastable supersymmetry breaking and multitrace deformations of SQCD,” arXiv:0812.3213 [hep-th].

[5] A. Giveon, D. Kutasov, J. McOrist and A. B. Royston, “D-Term Supersymmetry Breaking from Branes,” arXiv:0904.0459 [hep-th].

[6] M. J. Duncan and L. G. Jensen, “Exact tunneling solutions in scalar field theory,” Phys. Lett. B 291, 109 (1992).

[7] A. Amariti and M. Siani, “R-symmetry and supersymmetry breaking in 3D WZ models,” JHEP 0908, 055 (2009) [arXiv:0905.4725 [hep-th]].

[8] A. Giveon and D. Kutasov, “Seiberg Duality in Chern-Simons Theory,” Nucl. Phys. B 812, 1 (2009) [arXiv:0808.0360 [hep-th]].

[9] V. Niarchos, “Seiberg Duality in Chern-Simons Theories with Fundamental and Adjacent Matter,” JHEP 0811, 001 (2008) [arXiv:0808.2771 [hep-th]].

[10] V. Niarchos, “R-charges, Chiral Rings and RG Flows in Supersymmetric Chern-Simons-Matter Theories,” arXiv:0903.0433 [hep-th].

[11] A. Amariti, D. Forcella, L. Girardello and A. Mariotti, “3D Seiberg-like Dualities and M2 Branes,” arXiv:0903.3222 [hep-th].

[12] E. Witten, “Supersymmetric index of three-dimensional gauge theory,” arXiv:hep-th/9903005.

[13] A. Giveon and D. Kutasov, “Brane dynamics and gauge theory,” Rev. Mod. Phys. 71, 983 (1999) [arXiv:hep-th/9802067].

[14] T. Kitao, K. Ohta and N. Ohta, “Three-dimensional gauge dynamics from brane configurations with (p,q)-fivebrane,” Nucl. Phys. B 539, 79 (1999) [arXiv:hep-th/9808111].

[15] A. Giveon and D. Kutasov, “Gauge symmetry and supersymmetry breaking from intersecting branes,” Nucl. Phys. B 778, 129 (2007) [arXiv:hep-th/0703135].