Duality and Dynamical Supersymmetry Breaking
in Spin(10) with a Spinor

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We study $\mathcal{N} = 1$ supersymmetric Spin(10) chiral gauge theories with a single spinor representation and $N$ vector representations. We present a dual description in terms of an $\mathcal{N} = 1$ supersymmetric $SU(N - 5)$ chiral gauge theory with a symmetric tensor, one fundamental and $N$ antifundamental representations. The Spin(10) theory with $N = 0$ breaks supersymmetry at strong coupling; we study how this arises at weak coupling in the dual theory, which is a spontaneously broken gauge theory. Also, we recover various known dualities, find new dual pairs and generate new examples of dynamical supersymmetry breaking.

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1. Introduction

For supersymmetry to be relevant to the real world, it must be broken somehow. Dynamical supersymmetry breaking \[1\] is one interesting possibility. Until recently, few models with this feature had been found \[2,3\], but now a great many more are known \[4-7\]. In the last two years, there has been considerable progress in the understanding of the low-energy dynamics of supersymmetric field theories \[8\]. The unifying principle in this endeavor is duality \[9\]: two different gauge theories may have exactly the same long distance physics. But to establish duality, supersymmetry is essential.

In this letter, as well as in \[10,11\], the dynamical breaking of supersymmetry is studied in a theory which exhibits duality. We consider an \(\mathcal{N} = 1\) supersymmetric \(Spin(10)\) gauge theory. Its matter content consists of \(N\) superfields \(Q\) in the \(10\)-dimensional (vector) representation and a field \(P\) in the \(16\)-dimensional (spinor) representation of the gauge group and it has a tree-level superpotential \(W = 0\).

When \(N = 0\), the classical scalar potential has no flat direction. It was shown in \[3\] that supersymmetry is broken dynamically on the basis that no plausible low-energy description could be found satisfying the ’t Hooft anomaly matching conditions. In \[12\], the \(N = 1\) theory was used to prove that no vacuum state exists at weak coupling, consistent with the result of \[3\]. In this letter, we present new and independent evidence that no vacuum exists at strong coupling either.

To this end, we will study a dual for these \(Spin(10)\) theories. The dual “magnetic” theory we propose is, for \(N \geq 7\), an \(\mathcal{N} = 1\) supersymmetric gauge theory with gauge group \(SU(N - 5)\), with \(N\) matter superfields \(q\) in the antifundamental representation, a field \(q'\) in the fundamental representation, and a field \(s\) in the symmetric tensor representation. There are also elementary gauge singlets \(M\) and \(Y\) in the dual, to be identified by duality with composite gauge singlets \(M = Q^2\) and \(Y = QP^2\) of the \(Spin(10)\) theory. These fields interact via a superpotential \(W_{magnetic} = \text{det} s/\mu_1^{N-8} + MQsq/\mu_2^3 + Yqq'/\mu_3^2\). When \(N = 7\), an extra invariant term \(M^6Y^2\) must be added to \(W_{magnetic}\) in order to reproduce known results. The duality presented here is an extension of the one found in \[15\] for \(G_2\) and \(Spin(7)\) (see also \[16\] for an intermediate step to \(Spin(8)\).)

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1. \(SO(10)\) and \(Spin(10)\) have the same algebra, but only the latter has spinor representations. \(Spin(10)\) has two inequivalent complex spinor representations \(16\) and \(\bar{16}\); while \(16 \otimes \bar{16}\) contains a singlet, \(16 \otimes 16\) does not.

2. As emphasized in \[3,4\], the scales \(\mu\) are crucial for the details of the duality; we will, however, set \(\mu = 1\), as our purposes do not require keeping track of them.
2. Checks of the Duality

The first step to analyze the physics is to find the independent gauge-invariant chiral operators. In the “electric” \textit{Spin}(10) theory, there are mesons \(M_{ij} = Q_i Q_j\) which are symmetric in their flavor indices \(i, j = 1, \ldots, N\), and \(Y_i = Q_i PP\). \(M\) and \(Y\) are also present in the magnetic theory, but as elementary fields. There are also a number of baryons or exotics, totally antisymmetric in flavor, namely \(B = Q_5 P^2\), \(D_n = Q^{6+2n} W^{2-n}_\alpha\), and \(E_n = Q^{5+2n} P^2 W^{2-n}_\alpha\), \((n = 0, 1, 2)\), where \(W_\alpha\) is the gauge superfield strength.

In the magnetic theory, there are baryons \(\tilde{B} \equiv q^{N-5}\) contracted with one \((N-5)\)-index epsilon tensor, \(\tilde{D}_n = q^{N-6-2n} s^{N-6-n} q' W^n_\alpha\) and \(\tilde{E}_n = q^{N-5-2n} s^{N-5-n} \tilde{W}^n_\alpha\) \((n = 0, 1, 2)\), contracted with two \((N-5)\)-index epsilon tensors. The operators \(\det s, q' s^{N-6} q', qsq, qq'\) are redundant because of the equations of motion following from the superpotential.

There are many consistency checks on this duality. First, as required by duality, the global symmetry of the magnetic theory is the same as that of the electric theory, namely \(SU(N) \times U(1) \times U(1)_R\) (there is no discrete symmetry). Under this symmetry the matter fields \(Q\) and \(P\) of the electric theory transform as \((N, -1, 1 - \frac{8}{N+2})\) and \((1, \frac{1}{2} N, 1 - \frac{8}{N+2})\) while the matter fields of the magnetic theory \(s, q, q'\) transform as \((1, 0, \frac{2}{N-5}), (\tilde{N}, 1, \frac{8}{N+2} - \frac{1}{N-5})\) and \((1, -N, -1 + \frac{16}{N+2} + \frac{1}{N-5})\). Note that our choice of a basis for the \(R\)-symmetry \(U(1)_R\) is arbitrary. A highly non-trivial check is that the ’t Hooft anomaly matching conditions are satisfied at the origin of the moduli space of vacua. The symmetries, holomorphy and weak coupling forbid any dynamically generated superpotential for \(N \geq 6\), as is clear from the above choice of \(R\)-charges (we use the usual convention that the superpotential has \(R\)-charge 2.) The symmetries, holomorphy, and smoothness near the origin \(M, Y, s, q, q' = 0\) uniquely determine the magnetic superpotential for \(N > 7\).

There is a one-to-one correspondence, preserving all global symmetries, between the gauge-invariant operators \((M, Y, B, D_n, E_n)\) of the electric theory and the operators \((M, Y, \tilde{B}, \tilde{D}_n, \tilde{E}_n)\) of the magnetic theory. Some of these operators are constrained. In a similar duality mapping \[16\] we have shown that the mapping of the constraints is consistent, using the chiral anomaly and non-perturbative dynamics. However, we did not perform such checks in detail for the operators here.

\[\text{In } B \text{ the two } P \text{ fields are combined in a } 126 \text{ representation (a self-dual five-vector-index tensor) while in } E_n \text{ the two } P \text{ fields are combined in a vector representation; all of the vector indices are then contracted with a ten-index epsilon tensor.}\]
For \( N \geq 22 \), the electric theory is not asymptotically free, so it flows to a free theory of \( \text{Spin}(10) \) gluons and quarks in the infrared, as does the magnetic theory. For \( 7 \leq N \leq 21 \), we expect, but cannot prove, that the electric and magnetic theories are in a non-trivial interacting superconformal phase at the origin.

To check the duality further, we consider certain deformations. If we give an expectation value \( \langle Y_N \rangle \neq 0 \) to a component of \( Y \), the \( \text{Spin}(10) \) group is broken to \( \text{Spin}(7) \) by the Higgs mechanism. The fields \( Q_N, P \) are eaten by the massive gauge bosons and \( N - 1 \) eight-dimensional spinors \( \hat{Q} \) of \( \text{Spin}(7) \) remain massless. The other two components of \( Q_i, i = 1, \ldots, N - 1 \), become decoupled singlet fields. In the magnetic theory, the expectation value for \( Y_N \) gives mass to \( q_N \) and \( q'_N \), leaving an \( SU(N - 5) \) theory with \( N - 1 \) fields \( q \), a symmetric tensor \( s \), a superpotential \( W = \det s + [\langle M \rangle + M] qsq + Yqq' \). One direction is to explore these new theories in detail. Instead, for conciseness, we will mention only two of them here, and show that they match to previously known dualities.

Consider the case \( r = 1 \). For convenience, let us take \( N + 1 \) vectors \( Q_i, i = 0, 1, \ldots, N \) in the electric \( \text{Spin}(10) \) theory. Say \( \langle M_{00} \rangle \neq 0 \). The resulting \( \text{Spin}(9) \) theory has \( N \) vectors and a sixteen-component spinor \( \hat{P} \). While no mass term for \( P \) could be written in \( \text{Spin}(10) \), a mass term exists for \( \hat{P} \) in \( \text{Spin}(9) \), namely \( Y_0 = \langle Q_0 \rangle P^2 \) in terms of \( \text{Spin}(10) \) fields. When we add \( Y_0 \) to the superpotential of \( \text{Spin}(9) \), the electric theory becomes \( SO(9) \) with \( N \) vector representations \( \hat{Q} \). When we add \( Y_0 \) to the magnetic superpotential \( W = \det s + \langle M_{00} \rangle q^0 sq^0 + Mqsq + Yqq' + Y_0 \), the operator \( q^0 q' \) condenses. This breaks the dual \( SU(N - 4) \) gauge group to \( SU(N - 5) \). From the equation of motion for \( s \), we have \( s^{-1} \det s + Mqq = -\langle M_{00} \rangle q^0 q^0 \). Since \( M \) is kept arbitrary in the electric theory, this implies \( s^{-1} \det s \neq 0 \). The gauge group is thus broken further to \( SO(N - 5) \). It can be checked that, after all the massive fields are integrated out, the remaining fields are \( N \)
vectors $\hat{q}^i$ of $SO(N - 5)$ with the singlets $\hat{M}_{ij}, i, j = 1, \ldots, N$, some decoupled singlet fields and the superpotential $\hat{M}\hat{q}\hat{q}$, as expected from the duality between $SO(9)$ and $SO(N - 5)$ with $N$ vectors [13].

For our second example, we will derive one more new dual pair and use it to relate the $Spin(10)$ duality to the $Spin(8)$ duality of [16]. We start with a similar set-up as in the previous paragraph: the $Spin(10)$ theory with $N + 1$ vectors $Q$. First, add $Y_0$ to the superpotential; then go along the flat direction $\langle M_{0N} \rangle \neq 0$. In the electric theory, $Spin(10)$ is broken to $Spin(8)$. The 16-spinor $P$ splits into an $8_s$-spinor $\hat{P}$ and an $8_c$-conjugate-spinor. Note that under this breaking, the superpotential $\hat{P}$ and $W$ $SU$ are a singlet or a fundamental under $Spin(8)$; then go along the flat direction $\langle \rangle \neq 0$. In the magnetic theory, the case $N = 8$ gives a mass only to the $8_c$; the mass term for $\hat{P}$ is $Y_N$. The result is $Spin(8)$ with $N - 1$ vectors $\hat{Q}$, a spinor $\hat{P}$ and $W_{electric} = U \hat{P}\hat{P}$, where $U$ is the $Spin(8)$-singlet component of $Q_0$ which neither gets an expectation value nor is eaten by massive gauge bosons. In the magnetic $SU(N - 4)$ theory, $W = \det s + Mqsq + \langle M_{0N} \rangle q^0sq^N + Yqq' + Y_0$. As above, $SU(N - 4)$ is first higgsed to $SU(N - 5)$ by $\langle q^0q' \rangle \neq 0$. From the term $\langle M_{0N} \rangle (q^0) sq^N$, the components of $s$ which are a singlet or a fundamental under $SU(N - 5)$ are massive. Consequently, the det $s$ term is eliminated. The result of integrating out the massive fields is $W = \hat{M}_{ij}\hat{q}^i\hat{q}^j, i, j = 1, \ldots, N - 1, \alpha, \beta = 1, \ldots, N - 5$, as expected from [16].

We note that by adding singlets to both the electric and the magnetic theories, the duality for $Spin(10)$ as an electric theory is trivially converted into one for $SU(N - 5)$ as the electric theory with the same charged matter content as the above magnetic theory but no singlet fields. Specifically, the dual to $SU(N - 5)$ with fields $s, q, q'$ and superpotential $W = \det s$ is $Spin(10)$ with fields $Q, P$, singlets $\hat{M} = qsq, \hat{Y} = qq'$, and a superpotential $W = \hat{M}\hat{Q}^2 + \hat{Y}\hat{Q}\hat{P}^2$.

We now show how to flow to a duality in $\mathcal{N} = 2$ supersymmetric $SU(2)$ gauge theory found in [19]. Consider for the $Spin(10)$ gauge theory the case $N = 8$. Add to this theory seven gauge singlet fields $\hat{M}^8_j, j = 1, \ldots, 7$. Take the tree level superpotential to be $W_{electric} = Q_8\hat{P}^2 + \hat{M}^8_j Q_8 Q_j$. Go along a flat direction with $\langle M_{ij} \rangle \propto \delta_{ij}, i, j = 1, \ldots, 7$. The result after integrating out the massive fields is a $Spin(3) \cong SU(2)$ gauge theory with a triplet $\hat{Q}_8$ and eight doublets $\hat{P}$, and with the superpotential $\hat{Q}_8\hat{P}\hat{P}$ that renders the theory $\mathcal{N} = 2$ supersymmetric (and ultraviolet finite). Now let us study the dual. It is an $SU(3)$ gauge theory with the superpotential $W = Mqsq + Yqq' + \det s + Y_8 + \hat{M}^8_j M_{8j},$ with $\det s = s^3$. First $SU(3)$ is broken to $SU(2)$ by $\langle q^8q' \rangle$. The field $s$ decomposes into a triplet $\hat{s}$, a doublet $q^0$ and a singlet; the fields $M_{88}, Y, M_{8j}, \hat{M}^8_j$ and the $SU(2)$-singlet components of the charged fields are all massive and may be integrated out. Then, with
\[ \langle M_{ij} \rangle \propto \delta_{ij}, \ i, j = 1, \ldots, 7, \] the result is simply the superpotential \[ W = \hat{s}q^0\bar{q}^0 + \hat{s} \sum q^i\bar{q}^i. \] This magnetic theory also has \( \mathcal{N} = 2 \) supersymmetry and is isomorphic to the electric one \[ [19], \] as expected from Montonen-Olive duality \[ [20]. \]

4. Mass perturbations of the \( \text{Spin}(10) \) theory

We now add to the \( \text{Spin}(10) \) theory a mass term \[ W = mQQ = mM, \] where \( m \) has rank \( r \). We integrate out the massive vectors and flow to a \( \text{Spin}(10) \) theory with \( N - r \) vectors. In the dual, the operator \( qsq \) gets an expectation value of rank \( r \), breaking the color group from \( SU(N - 5) \) to \( SU(N - 5 - r) \). The duality is clearly preserved when \( r \leq N - 8 \). If \( r = N - 7 \), the magnetic gauge group is \( SU(2) \), the field \( s \) is massive and \[ W = M^0Y^2 + M^2(q^2)^2 + Yqq' \] (as noted above, the dynamical source of the \( M^6Y^2 \) term is unknown.) If \( r \geq N - 6 \), the \( SU(2) \) group is spontaneously broken. At this stage the electric theory confines. (In some regions of moduli space the electric gauge group is broken and the magnetic theory confines; but there is no invariant distinction between the confining and Higgs phases in this theory.) We can keep on studying the dual which is a theory of singlets. Alternatively, we can study the electric theory directly using the techniques of \[ [17]. \]

First consider \( r = N - 6 \), leaving six massless vectors. Both analyses yield the result that the light fields are the independent gauge singlets \( M_{ij}, Y_i, B^i \) (recall \( B^i \) is \( \epsilon^{i\overline{j}_1...\overline{j}_5}Q_{j_1}...Q_{j_5}P^2 \) and \( q^i \) in the electric and magnetic theories respectively) satisfying the two constraints

\[ \langle Y_iB^i \rangle = 0 \quad \text{and} \quad \langle M^5Y^2 + MB^2 \rangle = C \] (4.1)
on their expectation values, where \( C \) is a constant. The ’t Hooft anomaly matching conditions are satisfied. In the classical limit \( \Lambda_{\text{electric}} \to 0 \), we have \( C \to 0 \) and these constraints follow from \( \text{Spin}(10) \) Fierz identities. We may further check these constraints in the following two ways. Giving an expectation value \( \langle M \rangle \) of rank 6 for the fields \( M \) that have not been integrated out, we find that the gauge group is broken to \( \text{Spin}(4) \approx SU(2)_L \times SU(2)_R \); each gauge group has four doublets \( P_a^S \) \( (a = 1, 2, 3, 4; S = L, R) \), coming from the spinor \( P \), whose mesons \( V_{ab}^S = P_a^SP_b^S \) satisfy the two constraints \( \text{Pf } V_S = \Lambda^4_S \) \( (S = L, R) \) \[ [17]. \] The operators \( V \) are linear combinations of \( Y \) and \( B \) and the two \( \text{Spin}(10) \) constraints flow to linear combinations of the \( \text{Spin}(4) \) constraints. Similarly, along an \( SU(3) \) flat direction of the \( \text{Spin}(10) \) theory, where the massless field \( \langle M \rangle \) has
rank 2 and \( \langle Y \rangle \neq 0 \), three flavors of \( 3, \bar{3} \) remain. This theory satisfies only one constraint \([17]\). The two constraints of the Spin(10) theory relate a reducible \( SU(3) \) operator to irreducible ones, and can be combined into a single constraint on mesons and baryons of \( SU(3) \).

For \( r = N - 5 \), a superpotential proportional to \( 1/(M^4 Y^2 + B^2) \) is obtained. It is generated by instantons in the electric theory. When \( N - r = 1, 2, 3, 4 \) vectors remain massless, gaugino condensation in an unbroken subgroup of Spin(10) generates a superpotential proportional to \( [1/M^{N-r-1}Y^2]^{1/(6-N+r)} \). Clearly the equation of motion for \( Y \) requires supersymmetric solutions to be at infinite values of the fields.

5. Dynamical SUSY breaking

We now wish to study supersymmetry breaking. As mentioned in the introduction, it is believed that Spin(10) with one \( 16 \) and no \( 10 \)s breaks SUSY dynamically\( ^4 \). This theory has no chiral gauge-invariant operators and thus no flat directions; it is therefore very strongly coupled. Recently, Murayama \([12]\) studied the theory with \( N = 1 \). Classically there is a moduli space, which allowed him to study the theory at weak coupling. Under mass and Yukawa coupling perturbations, he showed that SUSY is broken at weak coupling in the Higgs phase. As the mass for the vector goes to infinity, the SUSY-breaking vacuum moves to the origin. There the theory is strongly coupled and theoretical control is lost.

Here, we present new evidence that SUSY is dynamically broken. We study the strong coupling regime of the theory of one spinor and \( N \) massive vectors, using the weakly coupled magnetic description. The dual provides the correct degrees of freedom near the origin of the moduli space: the fields \( M, Y, q \), etc. We can look for a supersymmetric vacuum by seeking solutions of the equations of motion for these fields. When small masses are given to all \( N \) vectors, these equations have no solution. By holomorphy, we may continue these masses to infinite values without restoring supersymmetry\( ^5 \), implying that SUSY is broken in the Spin(10) theory with one spinor and no vectors.

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\(^4\) We recall that a solution of the ’t Hooft anomaly matching conditions for unbroken \( U(1)_R \) was found in \([3]\), but was discarded because of its complexity. We see no trace of such a solution in our analysis of the dual theory.

\(^5\) Supersymmetry can be restored only on surfaces of complex codimension one in parameter space, so there exist curves along which all masses may be taken to infinity.
We will concentrate on the case $N = 6$. At this value of $N$ we have a whole moduli space of vacua where the 't Hooft anomaly matching conditions are satisfied, but no magnetic gauge dynamics which could introduce non-perturbative phenomena. The cases $1 \leq N \leq 5$, where the only vacua are at infinite field expectation values, can be derived from the case $N = 6$ (using holomorphy.) For $N \geq 7$, we may (by holomorphy) take all but six vectors to be very massive; the Higgs mechanism breaks the magnetic gauge group leaving the theory of singlets dual to the case $N = 6$.

For $N = 6$, the singlet fields satisfy the two constraints (4.1), which we implement in the superpotential using Lagrange multipliers $X_1, X_2$. Now add to the superpotential $mM + yY$; SUSY should be broken for generic values of the coupling $y$ and of the rank six matrix $m$. We claim the equations of motion for the magnetic superpotential $W = mM + yY + X_1(M^5Y^2 + MB^2 - C) + X_2YB$ have no solutions in its region of validity (near the origin of moduli space, at strong coupling.) From $M_{ij}(\partial W/\partial M_{ij}) = 0$ we learn finite $\langle M \rangle$ implies finite $X_1$. Noting that $(M^4Y^2)^{ij}Y_j = 0$ by antisymmetry, and using $Y_j(\partial W/\partial M_{ij}) = 0$ and $Y_jB^j = 0$, we find $m^{ij}Y_j = 0$; since $m$ is rank six, $Y = 0$. But $B^iB^j$ has rank at most one, so $\det[\partial W/\partial M] = \det m = 0$, contrary to assumption. This phenomenon is to be interpreted as dynamical supersymmetry breaking in the strongly coupled electric theory and as tree level, O’Raifeartaigh-type breaking in the infrared-free magnetic theory of singlets.

One may use other duality transformations to produce more theories that break SUSY. Consider the magnetic theory for $N = 7$, which is $SU(2) \approx Sp(1)$ with eight doublets $q$ and $q'$ and with $W = M^6Y^2 + M^2(q^2)^2 + Y qq' + mM$. We now dualize it according to the $Sp$ duality of [9]. We obtain a dual representation which is $Sp(1)$ with eight doublets $d$ and $d'$, and gauge singlets $M, Y, \hat{B} = q^2$ and $\hat{Y} = qq'$. Integrating out the massive fields $Y, \hat{Y}$ leaves the superpotential $W = M^6(dd')^2 + M^2\hat{B}^2 + \hat{B}dd + mM$. As expected, this theory breaks SUSY when $m$ has maximal rank 7. To see that, first note that if $m$ has rank 1, the field $\hat{B}$ may have rank 2, giving mass to two of the doublets $d$, leaving six doublets and causing the theory to confine [17]. (In other regions of moduli space the $Sp(1)$ gauge group is broken, but there is no distinction between confining and Higgs phases in this theory.) The spectrum and contraints of the resulting theory of singlets are then identical to the $N = 6$ theory studied in the previous paragraph, and the analysis for $m$ of higher rank follows immediately.

As another possibility, we can add two singlets $X, \tilde{X}$ (antisymmetric tensors of $SU(2)!)$ to the magnetic $SU(2)$ theory, with a potential $(X\tilde{X})^k$; this does not affect the dynamics.
This theory has a dual description \[18\] as $SU(4k + 2)$ with a flavor of fields $x, \bar{x}$ in the antisymmetric tensor representation, four flavors of fundamental and antifundamental representations, numerous singlets and a complicated superpotential. The specific dynamics of the resulting SUSY breaking depends on $k$. Similar tricks may be used to create huge classes of theories that break supersymmetry. Ideas of this type have been illustrated in \[10\].

It would be interesting to exhibit duality in a theory that flows down to the $SU(2)$ model with matter in the four-dimensional (spin 3/2) representation which was argued in \[8\] to break SUSY dynamically.

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References

[1] E. Witten, Nucl. Phys. B185 (1981) 513
[2] I. Affleck, M. Dine and N. Seiberg, Nucl. Phys. B256 (1985) 557
[3] I. Affleck, M. Dine and N. Seiberg, Phys. Lett. 140B (1984) 59
[4] A.E. Nelson and N. Seiberg, hep-ph/9309299, Nucl. Phys. B416 (1994) 46
[5] K. Intriligator, N. Seiberg and S. Shenker, hep-ph/9410203, Phys. Lett. 342B (1995) 152
[6] E. Poppitz and S. P. Trivedi, hep-th/9507169, EFI-95-44
[7] M. Dine, A.E. Nelson, Y. Nir and Y. Shirman, SCIPP 95/32, UW-PT/95-08, WIS-95/29/Jul-PH, hep-ph/9507378
[8] K. Intriligator and N. Seiberg, hep-th/9509060, RU-95-48, IASSNS-HEP-95/70
[9] N. Seiberg, hep-th/9411149, Nucl. Phys. B435 (1995) 129
[10] K. Intriligator and S. Thomas, SLAC-PUB-95-7041, to appear
[11] P. Pouliot, hep-th/9510148, RU-95-66
[12] H. Murayama, hep-th/9505082, Phys. Lett. 355B (1995) 187
[13] K. Intriligator and N. Seiberg, hep-th/9503179, Nucl. Phys. B444 (1995) 125; hep-th/9506084, RU-95-40, IASSNS-HEP-95/48
[14] D. Kutasov, A. Schwimmer and N. Seiberg, hep-th/9510222, EFI-95-68, WIS/95/27, RU-95-75
[15] P. Pouliot, hep-th/9507018, Phys. Lett. 359B (1995) 108
[16] P. Pouliot and M. Strassler, hep-th/9510228, RU-95-67
[17] N. Seiberg, hep-th/9402041, Phys. Rev. D49 (1994) 6857
[18] K. Intriligator, R.G. Leigh and M.J. Strassler, hep-th/9506118, RU-95-38
[19] N. Seiberg and E. Witten, hep-th/9407087, Nucl. Phys. B426 (1994) 19; hep-th/9408099, Nucl. Phys. B431 (1994) 484
[20] C. Montonen and D. Olive, Phys. Lett. 72B (1977) 117