Manifolds of Fixed Points and Duality in
Supersymmetric Gauge Theories

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There are many physically interesting superconformal gauge theories in four dimensions. In this talk I discuss a common phenomenon in these theories: the existence of continuous families of infrared fixed points. Well-known examples include finite $\mathcal{N} = 4$ and $\mathcal{N} = 2$ supersymmetric theories; many finite $\mathcal{N} = 1$ examples are known also. These theories are a subset of a much larger class, whose existence can easily be established and understood using the algebraic methods explained here. A relation between the $\mathcal{N} = 1$ duality of Seiberg and duality in finite $\mathcal{N} = 2$ theories is found using this approach, giving further evidence for the former. This talk is based on work with Robert Leigh (hep-th/9503121).

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I. INTRODUCTION

In this lecture, I will be presenting a new characterization of an important and common (but little known) phenomenon in four-dimensional $\mathcal{N} = 1$ supersymmetric gauge theories. This phenomenon, which goes by the name of “manifolds of fixed points” or “exactly marginal operators”, is familiar in two-dimensional field theory but has been little studied in four dimensions. A fixed point of a gauge theory is of course a point in the space of coupling parameters where the theory is conformally invariant; all coupling constants are truly constant, with all beta functions vanishing and all of the physics scale-invariant. A theory may have many isolated fixed points, but it may also possess a continuous set — a manifold — of fixed points. In this case, starting from one such point, the Lagrangian of the theory can be smoothly deformed in such a way that the quantum theory remains scale invariant. The operator which deforms the Lagrangian is said to be an exactly marginal operator; in four dimensions, its dimension is exactly four at every fixed point in the manifold.

In general, dimensions of operators in quantum theories cannot be computed exactly. However, in certain theories, including $\mathcal{N} = 1$ supersymmetric gauge theory in four dimensions, there is enough symmetry in the theory that certain anomalous dimensions are exactly known at fixed points, so operators can be shown to be exactly marginal and the existence of manifolds of fixed points can be confirmed. This fact was proven in Ref. [1], on which this talk is based.

The methods which are necessary for this purpose are remarkably simple; they merely require algebra. By this I do not mean abstract group theory — I mean the algebra we learned when we were children. From these methods we can make exact statements about the behavior of strongly coupled gauge theories, which allows us to clarify many old results and leads to many new ones. One application regards the relation between the electric-magnetic duality of $\mathcal{N} = 1$ theories [2] and that of $\mathcal{N} = 2$ finite theories [3,4]; it can be shown that the first follows from the second. [1]

Some of our results on finite models are related to earlier work of Lucchesi, Piguet and
II. RENORMALIZATION GROUP FLOW

Given a quantum field theory, which has gauge group $G$ and is weakly coupled in the ultraviolet, how does it change (flow) under scaling transformations? What is its behavior in the infrared (IR)? There are a number of possibilities, which include:

1. No light particles in the IR. The theory confines or is completely Higgsed, and there is a mass gap.

2. Only Goldstone bosons in the IR. The theory confines or is Higgsed, but global symmetries are broken spontaneously (as in QCD).

3. Weakly coupled gauge theory, with gauge group $H \subset G$, in the IR. The theory is either partly broken (as in the electroweak interactions) or partly confined (with light composite scalars and fermions potentially present.)

4. Weakly coupled gauge theory, with some other gauge group $G'$, in the IR. The gauge group $G$ is confined and a new gauge group emerges, with composite gauge bosons and matter emerging at low energies.

5. Interacting conformal fixed point in the IR. Many gauge theories may flow to the same such point.

Answers 4 and 5 are the most exotic, and only recently has it been realized [2,6] that these are in fact common situations in supersymmetric gauge theory.

III. THEORIES WITH FIXED POINTS.

I now turn to the specific issue of theories with IR fixed points. The fact that there are gauge theories in four dimensions that have non-trivial interacting fixed points has long been
known. One need only consider the beta function for QCD with $N_c$ colors and $N_f$ flavors of quarks in the fundamental representation:

$$\beta(g) = -\frac{g^3}{16\pi^2}\sum_0^\infty \left[\frac{g^2}{16\pi^2}\right]^k b_k,$$

$$b_0 = \frac{11}{3}N_c - \frac{2}{3}N_f; b_1 = \frac{N_c}{3}\left[34N_c - \left(13 - \frac{3}{N_c}\right)N_f\right].$$

(1)

Here $b_n$ is the $n$-loop contribution to the beta function. For asymptotically free theories, all the coefficients $b_k$ are of order at most $(N_c)^{k+1}$. If we take $N_f$ very close to but less than $\frac{11}{2}N_c$, for $N_c$ large (though not infinite), then the leading term in the beta function will be negative and of order 1 while the next-to-leading term will be positive and of order $g^2N_c^2$. The beta function therefore appears to have a zero near $g \sim 4\pi/N_c$. Can we trust this zero? Yes; for sufficiently large $N_c$ this value of the coupling is sufficiently small that the higher-order terms in the expansion of $\beta(g)$ will not change its basic functional form at small $g$. We may therefore conclude that the beta function is negative at very small $g^2$ but positive beyond some value. At the zero of the beta function lies a weakly interacting conformal field theory.

This same argument applies in supersymmetric QCD; for $N_f$ just less than $3N_c$ there are fixed points at weak coupling. In fact, Seiberg has argued \cite{2} (and by now the evidence is overwhelming \cite{3,4}) that such interacting field theories exist even for $N_f$ much less than $3N_c$, with couplings which in general are non-perturbatively strong.

Another theory which has conformal fixed points is $\mathcal{N} = 4$ supersymmetric gauge theory. In fact, this theory is conformal for every value of the gauge coupling. Let us consider a slightly more general theory; take an $\mathcal{N} = 1$ supersymmetric gauge theory, with three chiral superfields $\phi_1, \phi_2, \phi_3$ in the adjoint representation of the group, and with a superpotential

$$W = h f_{abc} \phi_1^a \phi_2^b \phi_3^c$$

(2)

where we do not assume $h$ and the gauge coupling $g$ are equal. In the space of all possible couplings $(g, h)$ there are non-trivial beta functions $\beta_g(g, h)$ and $\beta_h(g, h)$; but on the line $h = g$ both beta functions vanish by $\mathcal{N} = 4$ supersymmetry. Thus, in this $\mathcal{N} = 1$ theory
there is a line of fixed points, labelled by $g$; it can easily be shown that these fixed points are infrared stable. (In fact, for $SU(N)$ there is a much larger manifold of fixed points which contains $\mathcal{N} = 4$ as a subspace; see Ref. [1].) The variation of the Lagrangian with respect to $g$, $\frac{dC}{dg}$, is an exactly marginal operator on the line $h = g$.

![Diagram showing fixed points on a line](image)

**FIG. 1.** The $\mathcal{N} = 4$ superconformal theories lie on a line inside the $\mathcal{N} = 1$ theory (2). Arrows indicate renormalization group flow toward the infrared.

The reason for viewing the $\mathcal{N} = 4$ theory as a line inside the space of couplings of an $\mathcal{N} = 1$ theory is the following: there are many examples, known since the 1980s, in which an $\mathcal{N} = 1$ theory has a curve of fixed points inside the space of couplings, even though there is no special symmetry on this curve. From this point of view, $\mathcal{N} = 4$ is a special case of a much broader class of $\mathcal{N} = 1$ theories which have neither beta functions nor anomalous dimensions for their chiral superfields; their effective actions are finite. [7–9] These theories were the motivation for the Principle of Reduction of Couplings [10] mentioned in the talk of Dr. Kubo.

I will now show that there is a simple way to understand why these theories exist, and that there are many theories which are *not* finite which nonetheless have manifolds of fixed points.
IV. PROVING AN OPERATOR IS MARGINAL

The question I will now address is this: given a theory which is at a fixed point, whose Lagrangian is $\mathcal{L}$, and which has an operator $\mathcal{O}$ which is marginal (dimension-four) at the fixed point, under what circumstances is the theory with Lagrangian $\mathcal{L} + \epsilon \mathcal{O}$ still at a fixed point? This is equivalent to asking whether the dimension of operator $\mathcal{O}$ is still four when $\epsilon$ is non-zero. If the answer to this question is yes, then the operator is exactly marginal.

In general, this is a very difficult question to answer, even in two dimensions, unless the theory is exactly soluble. But remarkably, it is easy to answer this question for certain operators in $\mathcal{N} = 1$ $d = 4$ supersymmetry. This follows from three important and little known facts about these theories:

1. Anomalous dimensions $\gamma$ of charged matter fields are gauge invariant.

2. If a term $h\phi_1 \ldots \phi_k$ appears in the superpotential, then
   \[
   \beta_h \propto \left[ (k - 3) + \frac{1}{2} \sum_{i=1}^{k} \gamma_i \right]
   \] (3)
   where $\gamma_i$ is the anomalous mass dimension of the field $\phi_i$.

3. For the gauge coupling $g$, the beta function, in terms of the second Casimir invariant of the gauge group $C_2(G)$ and the index $l_R = 2T(R)$ of the representation $R$, is
   \[
   \beta_g \propto -\left( 3C_2(G) - \sum_i T(R_i) \right) + \sum_i T(R_i) \gamma_i \] (4)
   where the sums are over all matter fields. Notice the one-loop beta function coefficient $b_0 = 3C_2(G) - \sum_i T(R_i)$ appears in this formula.

The expression (3) follows directly from the non-perturbative non-renormalization theorem [11] for the superpotential, while (4) was proven to all orders in perturbation theory by Shifman and Vainshtein [12] (see also papers with Novikov and Zakharov [13]). The latter is believed also to hold non-perturbatively because the zeroes of the beta function are related,
by properties of the superconformal algebra of $\mathcal{N} = 1$ supersymmetry, \cite{14,15} to certain chiral anomalies, which can be computed at one loop.

What is critical to note about these beta functions is that both are linear functionals of the anomalous dimensions $\gamma_i$ of the matter fields $\phi_i$. These anomalous dimensions are arbitrarily complicated functions of the coupling constants, but the relations between beta functions and anomalous dimensions are very simple. We will now see how this leads to a powerful and elegant result.

The criterion for the existence of a fixed point is that all beta functions must vanish simultaneously. The equations $\beta_i = 0$ put $n$ conditions on $n$ couplings. Generically, we would expect these conditions are satisfied at isolated points in the space of couplings. However, because the beta functions are linear in the anomalous dimensions, it may happen that some of them are linearly dependent. If only $p$ of the beta functions are linearly independent, then the equations $\beta_i = 0$ put only $p$ conditions on the $n$ couplings, so we expect that solutions to the equations generically occur on $n - p$ dimensional subspaces of the space of couplings, leading to $n - p$ dimensional manifolds of fixed points.

Thus, if some of the beta functions are linearly dependent, and there is at least one generic fixed point somewhere in the space of couplings, then the theory will have a manifold of fixed points. Of course, it is possible that a theory has no fixed points at all (other than the free fixed point, which is not generic.) But in many theories, non-trivial fixed points are known or believed to exist, and in some cases these extend to entire manifolds.

V. EXAMPLES

As a first example of such a theory, consider $SU(3)$ with fields $Q^r, \tilde{Q}_s (r, s = 1, \ldots, 9)$ in the $3$ and $\bar{3}$ representations, with a superpotential

$$W = h \left( Q^1 Q^2 Q^3 + Q^4 Q^5 Q^6 + Q^7 Q^8 Q^9 + \tilde{Q}_1 \tilde{Q}_2 \tilde{Q}_3 + \tilde{Q}_4 \tilde{Q}_5 \tilde{Q}_6 + \tilde{Q}_7 \tilde{Q}_8 \tilde{Q}_9 \right).$$

The special properties of this theory were first noted in the 1980s. \cite{7,8} Notice that the one-loop gauge beta function is zero and that the superpotential preserves enough of the
flavor symmetry that all $Q^r$ and $\tilde{Q}_s$ have the same anomalous dimension $\gamma(g, h)$. The beta functions for the Yukawa coupling and gauge coupling are

$$\beta_h \propto \frac{3}{2} h \gamma(g, h); \quad \beta_g \propto -9 \gamma(g, h).$$

For both of these to vanish requires only the one condition $\gamma(g, h) = 0$. Since $\gamma(0, 0) = 0$, $\gamma(g, 0) < 0$ and $\gamma(0, h) > 0$ for small $g, h$, we conclude there is a line of fixed points passing through the origin of the space of couplings $(g, h)$.

![Diagram](image.png)

**FIG. 2.** The renormalization group flow near the fixed curve associated with the superpotential (5), shown schematically, in the plane of the gauge coupling $g$ and the Yukawa coupling $h$.

There are a number of important facts worthy of note:

1. On the fixed line, by definition, the theory is superconformal. Since all beta functions and anomalous dimensions vanish, the effective action on the line is finite (just as $\mathcal{N} = 4$ theories are finite.)

2. We have not determined the function $\gamma(g, h)$, but we didn’t need to. The existence and properties of the fixed line only required that $\gamma(g, h)$ be zero somewhere. Furthermore, if $g, h$ are small, we can compute $\gamma(g, h)$ in perturbation theory, while if $g, h$ are large this function is not useful anyway.
3. The line of fixed points is IR stable; if the couplings are near but off the line, they will
approach the line at low energy. Therefore, no fine tuning is required to put the theory
on the line, which means these superconformal theories are physically interesting.

There are an enormous number of models with similar manifolds of fixed points passing
through zero coupling; these can be and were identified in perturbation theory.

As another example, consider an $\mathcal{N} = 1$ supersymmetric $SU(N)$ gauge theory with fields
$Q^r, \bar{Q}_s$ ($r, s = 1, \ldots, N_f$) in the fundamental and antifundamental representations, a single
field $\Phi$ in the adjoint representation, and a superpotential

$$W = h\Phi Q^r \bar{Q}_r$$ (7)

where the repeated index is summed over. For $h = g$ this theory has $\mathcal{N} = 2$ extended
supersymmetry. In the case $N_f = 2N$, the theory is finite on the line $h = g$; we can see
this line of fixed points must exist using the fact that both $\beta_g$ and $\beta_h$ are proportional to
$\gamma_\Phi + 2\gamma_Q$ (by symmetry, all $Q^r, \bar{Q}_s$ have the same anomalous dimension). Again the line of
fixed points is infrared stable.

But now let us add a mass for the field $\Phi$ and integrate it out. The low energy theory is an
$\mathcal{N} = 1$ $SU(N)$ gauge theory with the $N_f = 2N$ fields $Q^r, \bar{Q}_s$ coupled by a non-renormalizable
superpotential

$$W = \frac{1}{2} m\Phi^2 + h\Phi Q^r \bar{Q}_r \rightarrow W_L = \frac{1}{2} y \left[ (Q^r \bar{Q}_s)(Q^s \bar{Q}_r) - \frac{1}{N}(Q^r \bar{Q}_r)(Q^s \bar{Q}_s) \right]$$ (8)

where an $SU(N)$ Fierz identity has been used to write the low energy superpotential $W_L$ in
terms of gauge-invariant bilinears. This superpotential preserves enough flavor symmetry
that all $Q^r, \bar{Q}_s$ have the same anomalous dimension $\gamma(g, y)$. The beta function for $y$ is
$\beta_y \propto 1 + 2\gamma$ while that of the gauge coupling is $\beta_g \propto (3N - N_f) + N_f \gamma$, which for $N_f = 2N$
is proportional to $\beta_y$. We therefore have a fixed point whenever $\gamma(g, y) = -\frac{1}{2}$.

When the couplings $g, y$ are small, $\gamma \sim Ay^2 - Bg^2$ for positive coefficients $A, B$; $|\gamma| \ll \frac{1}{2}$
at weak coupling, and any fixed point is therefore non-perturbative. The work of Seiberg strongly suggests that there is a value of $g = g_*$ for which the theory with $W_L = 0$ is at a
stable fixed point. If this is true, then we expect that \( \gamma(g, 0) > -\frac{1}{2} \) [\( \gamma(g, 0) < -\frac{1}{2} \)] for \( g < g^* \) [\( g > g^* \)]. There must therefore be a line of stable fixed points, passing through the Seiberg fixed point \((g, y) = (g^*, 0)\), which all have \( \gamma(g, y) = -\frac{1}{2} \).

![Diagram](image)

**FIG. 3.** The renormalization group flow near the fixed curve associated with the superpotential (8), shown schematically, in the plane of the gauge coupling \( g \) and the Yukawa coupling \( h \).

1. The fact that this theory has a quartic, non-renormalizable superpotential poses no difficulty. The low-energy theory is an effective theory valid below the mass of \( \Phi \). Above this mass the theory is renormalizable. In fact, one may redo the analysis in the renormalizable theory and obtain the same result concerning the existence of a fixed line.

2. The coupling \( y \) in (8) has dimension of inverse mass, so one might ask how it can appear in a superconformal field theory. The reason is that the anomalous dimension of the operator \((Q\bar{Q})^2\) is \( \frac{1}{2} \cdot 4\gamma = -1 \) at the fixed point; this exactly cancels the naive dimension of the coupling \( y \).

3. As before, the fixed line is IR stable, so no fine-tuning is required for the theory to be superconformal in the IR.

Again, the number of similar examples is enormous.
VI. APPLICATION TO $\mathcal{N} = 1$ DUALITY

While there is insufficient space to present the full relation between $\mathcal{N} = 2$ and $\mathcal{N} = 1$ duality, I will sketch the main ideas; further discussion is to be found in Ref. [1]. (In intervening months, an improved understanding of this relation has been achieved; more up-to-date discussion can be found in Ref. [16].)

The $\mathcal{N} = 2$ finite $SU(N)$ theory, whose line of fixed points was discussed above, has the property that at every point along the $\mathcal{N} = 2$ fixed line, there are (at least) two equivalent descriptions of the superconformal theory. We may refer to the two sets of variables as electric and magnetic. In both variables the model is a finite $\mathcal{N} = 2$ $SU(N)$ theory, but the magnetic coupling $\tilde{g}$ is the inverse of the electric coupling $g$, and the flavor quantum numbers are assigned differently in the two descriptions.

When a mass is given to the field $\Phi$ (this occurs similarly in both sets of variables) both descriptions flow to a low-energy theory consisting of an $\mathcal{N} = 1$ $SU(N)$ gauge theory with $N_f = 2N$, as above. Both theories have a superpotential of the form (8), but the coupling $\tilde{y}$ in the magnetic description is the inverse of $y$ in the electric description. If we take the coupling $g$ of the $\mathcal{N} = 2$ electric theory to be arbitrarily small, then the electric $y$ will be arbitrarily small as well; in the limit, the $\mathcal{N} = 1$ electric superpotential is zero, but the theory will still flow from weak coupling ($g = 0$) to its fixed point ($g = g_*$) since it is asymptotically free.

On the other hand, for large $\tilde{g}$ the magnetic theory will be at strong coupling, and will have a superpotential with a huge coefficient $\tilde{y}$. But we may introduce an auxiliary gauge singlet field $M^r_s$ to rewrite it as follows.

$$W_{mag} = -\frac{1}{2} \tilde{y} \left[ (q^r \tilde{q}_s)(q^s \tilde{q}_r) - \frac{1}{N} (q^r \tilde{q}_r)(q^s \tilde{q}_s) \right] \rightarrow M^r_s q_r \tilde{q}^s + \frac{1}{2\tilde{y}} (M^r_s M^s_r - \frac{1}{N} M^r_s M^s_r) . \quad (9)$$

In the limit that the magnetic theory is infinitely coupled (which is the same limit we

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1This section was omitted from the talk due to time constraints.
considered above for the electric theory) the mass of the meson $M$ is zero and we are left with the superpotential

$$W^\text{mag}_L = M^r_s q_r \tilde{q}^s. \quad (10)$$

We expect the low-energy $\mathcal{N} = 1$ magnetic theory to flow from infinite coupling back to its stable IR fixed point $\tilde{g} = \tilde{g}_s$.

We therefore conclude that the duality of the $\mathcal{N} = 2$ theory implies that an $\mathcal{N} = 1$ $SU(N)$ gauge theory with $N_f = 2N$ and superpotential $W = 0$ is dual to another $SU(N)$ theory with $N_f = 2N$ and mesons $M^r_s$ coupled by the superpotential $W = M q \tilde{q}$. This is exactly what was found by Seiberg in Ref. [2].

More details of this relation may be checked. [1] In particular, the mapping between operators found in Ref. [2] is also recovered. Using the mapping, one can flow away from the theory with $N_f = 2N$ and show that the $SU(N)$ theory with $N_f$ flavors has a magnetic dual which has $SU(N_f - N)$ gauge group, $N_f$ flavors, singlets $M^r_s$, and the superpotential (10). [1]

\textbf{VII. CONCLUSIONS}

1. Manifolds of fixed points are a common phenomenon in $\mathcal{N} = 1$ supersymmetric gauge theories.

2. These manifolds can easily be found, using simple algebraic methods.

3. The “finite” $\mathcal{N} = 1$ models discovered ten years ago are easily understood using this language. Proving a model is finite is often trivial with these methods.

4. Many “non-finite”, non-renormalizable, strongly coupled manifolds of fixed points have also been uncovered using this approach.

5. These fixed points are usually IR stable, and are therefore of physical interest.
6. Manifolds of fixed points appear to play an important role in duality, as manifested by the relation discussed above between duality in finite $\mathcal{N} = 2$ models and Seiberg’s $\mathcal{N} = 1$ duality.

7. Beyond those involving the Principle of Reduction of Couplings [10], no applications to phenomenology have yet been suggested, but approximate fixed points have the potential to play an important role in supersymmetric model building.

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