On the Vacuum Structure of the 3-2 Model

Tomer Shacham

Racah Institute of Physics,
The Hebrew University,
Jerusalem 91904, Israel

E-mail: Tomer.Shacham@phys.huji.ac.il

Abstract: The 3-2 model of dynamical supersymmetry breaking is revisited, with some incidentally new observations on the vacuum structure. Extra matter is then added, and the vacuum structure is further studied. The parametric dependence of the location of the vacuum provides a consistency check of Seiberg duality.

Keywords: Dynamical Supersymmetry Breaking, Vacuum Structure, Seiberg Duality
1 Introduction

The first step towards understanding the spectrum and dynamics of a quantum field theory is to study its vacuum structure. In the context of asymptotically free gauge theories, it is considered calculable if the entire gauge group is higgsed by expectation values parametrically larger than the scale of strong coupling. When this happens, quantum corrections are relatively small and a classical analysis of the vacuum structure is justified.

When the vacuum is strongly coupled, there is currently no description of its structure. However, in certain cases, $\mathcal{N} = 1$ supersymmetry (SUSY) allows an insight to the low energy physics given by Seiberg duality \cite{1}. Since this is a strong-weak duality, the low energy descriptions of some asymptotically free theories are IR free. The classical analysis of the vacuum structure is then reliable for small values of the charged fields.

If SUSY is a symmetry of nature, it must be broken. For reasons of naturalness, it should be broken dynamically. The 3-2 model \cite{2} (recently reviewed in \cite{3}) is an example of dynamical SUSY breaking. Due to its simplicity (or perhaps the resemblance to the Standard Model) it has been extensively used as a model building tool \cite{4, 5}.

In this note, the vacuum structure of the 3-2 model is revisited. A controllable vacuum is shown to exist in a large region of parameter space. Equipping the model with extra matter allows a low energy description to be provided by the Seiberg dual theory. The dependence of the vacuum structure on the parameter space of the theory provides a consistency check of the duality.

Since the duality still lacks a rigorous proof, consistency checks are of theoretical value. Available evidence includes the following. The dual theories have the same non-anomalous global symmetries; anomaly matching conditions for these symmetries are satisfied. If a SUSY vacuum exists, the dimension of the moduli space is the same in both theories.
Other nontrivial checks include corresponding dualities in $\mathcal{N} = 2$ theories [6], matching of the superconformal index [7, 8] and a brane construction in type IIA string theory [9, 10].

Following the footsteps of [11], new non-holomorphic evidence is presented here. The main idea is that the electric and magnetic theories cannot both be valid in overlapping regions of parameter space, as it is inconceivable that there are two different weakly coupled descriptions of the same physics. It should be noted that since the 3-2 model does not have a SUSY vacuum for any choice of parameters, evidence of the duality collected here is independent of previous checks.

This note is organized as follows. Section 2 includes a brief review of the 3-2 model and an analysis of its vacuum structure. In section 3 extra matter is added to the model in order for a dual description to exist and the vacuum structure is further studied.

2 The 3-2 Model

The 3-2 model is an $\mathcal{N} = 1$ supersymmetric theory with gauge group $SU(3)_C \times SU(2)_L$ and matter content

| $[SU(3)_C]$ | $[SU(2)_L]$ |
|-------------|-------------|
| $Q_A^\alpha$ | $\Box$ | $\Box$ |
| $\tilde{Q}_a^\alpha$ | $\Box$ | $1$ |
| $L_A$ | $1$ | $\Box$ |

where $\alpha = 1, 2$ is a flavor index. In the limit $\Lambda_3 \gg \Lambda_2$, the vacuum structure can be understood as follows. At first, consider $SU(2)_L$ to be a global symmetry. A non-perturbative superpotential

$$W_{ADS} = \frac{2\Lambda^7}{\det (QQ)}$$

(2.1)

pushes the fields away from the origin of field space. It is therefore natural to describe the vacuum in terms of the microscopic fields rather than the gauge invariants, to allow for a canonical Kähler potential at large fields. In this regime, $W_{ADS}$ is generated by an $SU(3)$ instanton. In attempt to stabilize this runaway, one may add interactions at tree level:\footnote{This superpotential is generic; a term $\tilde{\lambda} \tilde{Q}^2 QL$ could be removed by a redefinition of $\tilde{Q}$.}

$$W_{tree} = \lambda \tilde{Q}^1 Q L.$$ 

(2.2)

To a first approximation, the vacuum can be found on the $SU(3)_C$ D-flat directions as this gauge coupling is the largest parameter in the theory. These are given by

$$Q_{a}^{A\dagger} \tilde{Q}_{a}^{b} - \tilde{Q}_{a}^{\alpha} \tilde{Q}_{a}^{\dagger \alpha} = 0.$$ 

(2.3)
Since $\mathcal{W}_{\text{tree}}$ breaks the flavor symmetry in the $\bar{Q}$ sector, the most general form of $\bar{Q}$ which respects (2.3) is $Q^T G$ with $G \in SU(2)$. A convenient parameterization is

$$ G = \begin{pmatrix} e^{i\phi} \cos \theta & e^{i\chi} \sin \theta \\ -e^{-i\chi} \sin \theta & e^{-i\phi} \cos \theta \end{pmatrix}. \quad (2.4) $$

Modulo gauge transformations, $Q^a_3$ has the form $\begin{pmatrix} a & 0 & 0 \\ 0 & b & 0 \end{pmatrix}^T$ with $a, b \in \mathbb{R}$. After scaling all fields in “units” of $\lambda^{-1/7} \Lambda$, the F-term is $\mathcal{F} \Lambda^4 \lambda^{10/7}$ where

$$ \mathcal{F} \equiv \frac{4}{a^4 b^6} + b^2 l_1^2 + \cos^2 \theta \left( 4 \left( \frac{a^2 + b^2}{a^6 b^4} \right) + a^2 \left( a^2 + l_1^2 + l_2^2 \right) \right) - \frac{8}{a^2 b^2} l_1 \cos (\theta_1 - \chi) \sin \theta + $$

$$ + (a, b, l_1, l_2, \theta_1 - \chi, \theta) \rightarrow (b, a, l_2, l_1, \theta_2 + \phi, \theta + \pi/2) \quad (2.5) $$

and $L_A$ is parametrized as $(l_1 e^{i\theta_1}, l_2 e^{i\theta_2})$. Note that $\mathcal{F}$ depends only on the combinations $(\theta_2 + \phi)$ and $(\theta_1 - \chi)$. This happens because the scalar potential has additional Abelian global symmetries; these can be used to set $\phi = \chi = 0$. One can check that $\mathcal{F}$ has remaining runaways. These are lifted by the weak gauging of $SU(2)_L$, which generates a D-term $\mathcal{D} \Delta \Lambda^4 \lambda^{10/7}$ where

$$ \mathcal{D} \equiv \frac{1}{8} \left( 4l_1^2 l_2^2 + (a^2 - b^2 + l_1^2 - l_2^2)^2 \right) \quad (2.6) $$

and $\Delta \equiv (g_2/\lambda)^2$. The full scalar potential is then $\Lambda^4 \lambda^{10/7} (\mathcal{F} + \Delta \mathcal{D})$.

The original analysis was carried out in the limit $\Delta \gg 1$, where the vacuum is found at the minimum of $\mathcal{F}|_{\mathcal{D}=0}$. The VEVs of all fields scale like $\lambda^{-1/7} \Lambda$ and the vacuum energy scales like $\lambda^{10/7} \Lambda^4$. The smallness of $\lambda$ ensures both that the gauge group is higgsed at scales much higher than $\Lambda$ and small SUSY breaking.

Recently [11], the vacuum structure in the opposite limit was found. When $\Delta \ll 1$, it can be understood as follows. Consider $SU(2)_L$, a global symmetry and assume that $L_2$ is stabilized far from the origin. This generates a large mass $|\lambda L_2|$ for $Q_1, \bar{Q}^1$ and these fields decouple. For large values of the remaining flavor $Q_2, \bar{Q}^2$, the gauge group is spontaneously broken to $SU(2)_C$. Moduli space dependence of the gaugino condensate results in a runaway potential; this direction is lifted by weakly gauging $SU(2)_L$. The vacuum sits at $L_2 \sim Q_2 \sim \Lambda (\lambda \Delta)^{-1/7}$ and the vacuum energy scales like $\Lambda^4 \lambda^{10/7} \Delta^{3/7}$. The smallness of $\Delta$ justifies the integration out of the first flavor; the condition $\lambda L_2 \gg \Lambda$ is equivalent to $g_2 \ll \lambda^4$.

A new observation is that the non-perturbative superpotential (2.1) is generated by an instanton even in the regime $\Delta \ll 1$ where the D-term and F-term are comparable. As before, consider $SU(2)_L$ to be global.

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$SU(2)_L$ generators are normalized such that $\text{Tr} T^a T^b = \frac{1}{2} \delta^{ab}$. 
The potential is minimized at $\theta = 0$ for all values of $\Delta$. In the absence of a D-term, $\mathcal{F}$ has a runaway to
\[
\langle L \rangle = (0, 2a^{-4}b^{-2}), \quad \langle b \rangle \to \infty.
\tag{2.7}
\]
For $\Delta \ll 1$, the vacuum is found at the minimum of $(\mathcal{F} + \Delta \mathcal{D}) \big|_{\text{runaway}}$. On the runaway,
\[
\mathcal{F} \big|_{\text{runaway}} = a^4 + \frac{8}{a^4b^6}, \quad \mathcal{D} \big|_{\text{runaway}} = \frac{1}{8} \left( a^2 - b^2 - \frac{4}{a^8b^4} \right)^2.
\tag{2.8}
\]
Assume that $a \ll b^{-2/5}$ and so the term $a^2$ in $\mathcal{D}$ can be neglected, this will be justified by self-consistency. Rescale the fields (again) in order to incorporate $\Delta$ dependence:
\[
b \to b\Delta^{-1/7}, \quad a \to a\Delta^{3/28}.
\tag{2.9}
\]
Under this scaling, the full scalar potential is
\[
V = \Lambda^4 \lambda^{10/7} \Delta^{3/7} \left( a^4 + \frac{8}{a^4b^6} + \frac{1}{8} \left( b^2 + \frac{4}{a^8b^4} \right)^2 \right).
\tag{2.10}
\]
The vacuum energy is $E \approx 2.7 \Lambda^4 \lambda^{10/7} \Delta^{3/7}$ and the VEVs are given by
\[
\langle b \rangle \approx 1.6 \left( \lambda \Delta \right)^{-1/7} \Lambda, \quad \langle a \rangle \approx 1.0 \left( \lambda / \Delta^{3/4} \right)^{-1/7} \Lambda
\tag{2.11}
\]
with $\lambda$, $\Delta$ and $\Lambda$ reintroduced. This vacuum can be trusted as long as $\lambda^{5/3} \ll g_2 \ll \lambda$.

For $\Delta \sim 1$, a numerical calculation shows a smooth interpolation between the two regimes, presented in Fig. 1. In the window $\lambda^4 \lesssim g_2 \lesssim \lambda^{5/3}$, quantum corrections are not under control and the vacuum structure is unknown.

\footnote{A field configuration corresponding to $\theta = \pi/2$ is connected by a gauge transformation.}
As expected by the theorem proved in [11], the ratio $\Delta D/F$ is always bounded. It decreases monotonically as a function of $\Delta$:

$$\frac{\Delta D}{F} \sim \begin{cases} 
\frac{3}{4} & \Delta \to 0 \\
\Delta^{-1} & \Delta \to \infty.
\end{cases}$$ (2.12)

3 The 3-2 Model with Extra Matter

We now equip the 3-2 Model with two massive vector-like pairs of $SU(3)_C$, $\Psi^i, \bar{\Psi}^i$, in order for a weakly coupled dual theory to exist in some region of parameter space. We restrict to the limit $\Delta \gg 1$.

For very large masses, the vacuum structure does not know about the extra matter, and the low energy theory is essentially the same one previously discussed. However, the effective strong coupling scale is higher since the $\Psi$'s have been integrated out at the mass scale $m$; the scale matching is given by $\Lambda_{\text{eff}}^7 = \Lambda^5 m^2$. The vacuum is located at

$$a \sim b \sim L_2 \sim \left(\frac{\delta}{\lambda}\right)^{1/7} \Lambda$$ (3.1)

with $\delta \equiv (m/\Lambda)^2$. Note that this analysis of the vacuum structure is rigorous as long as $\delta \gg 1$.

When $m$ is lowered beneath $\Lambda$, there is no longer a Wilsonian sense to integrate out the $\Psi$'s. Using holomorphy, one concludes that the only possible superpotential is

$$W = \frac{2\Lambda^5 m^2}{\det (QQ)} + \lambda Q^j Q L + m \text{Tr} (\Psi \bar{\Psi}) .$$ (3.2)

Since there are no mixed terms involving both $\Psi$'s and $Q, \bar{Q}, L$, the vacuum remains at (3.1) and can be trusted for $\delta < 1$ in the regime $\delta \gg \lambda$ where it is controllable.

As the $\Psi$ mass is further reduced, the vacuum moves to strong coupling and a different analysis is required. If a vacuum exists in the small fields regime, the description of the low energy physics is given by the Seiberg dual “magnetic” theory. In the limit $\Lambda_3 \gg \Lambda_2$, only the $SU(3)_C$ sector is dualized. The magnetic theory consists of 4 pairs of “quarks” $q^i$ and “antiquarks” $\bar{q}_j$, $4 \times 4$ “mesons” $M^j_i$ and the doublet $L$. It is convenient to decompose the mesons and quarks as

$$M = \begin{pmatrix} \Phi^i_j & \Upsilon^i_j \\
\bar{Y}^j_A & Z^j_A \end{pmatrix} = \frac{1}{\Lambda} \begin{pmatrix} \Psi_i \bar{\Psi}^j & \Psi_i \bar{Q}^j \\
Q^j_A \bar{\Psi}^i & Q^j_A \bar{Q}^j \end{pmatrix}, \quad q = (x, \sigma^A), \quad \bar{q} = (\bar{x}, \bar{\sigma}_j).$$ (3.3)
The magnetic gauge group consists only of the untouched $SU(2)_L; \tilde{\Upsilon}^j, Z^j, \sigma$ and $L$ transform in their fundamental representation. Note that $\sigma$ is reminiscent of the baryon $\frac{1}{2} \Psi^i_a \Psi^j_b Q^k_A$.

The superpotential is

$$W = q^i M^j_i \bar{q}_j - \mu^2 \text{Tr} \Phi + \lambda \Lambda Z^1 L$$

(3.4)

with $\mu^2 \sim m\Lambda = \sqrt{\delta \Lambda^2}$.\(^6\)

A significant shortcut can now be taken. Consider a theory with the same field content but with $\lambda$ sent to zero and $SU(2)_L$ regarded as a global symmetry. The superpotential is given by

$$W' = \sigma^A Z^j_A \bar{\sigma}_j + \sigma^A \tilde{\Upsilon}^j_A \bar{x}_j + x \Upsilon \bar{\sigma} + x \Phi \bar{x} - \mu^2 \text{Tr} \Phi.$$  

(3.5)

This theory has been studied in [12], the relevant details are presented here for completeness. The F-term for $\Phi^j_i$ breaks SUSY by the rank condition. Minimization of the tree level potential is achieved by saturating one component of $x$

$$\langle x \rangle = \langle \bar{x} \rangle^T = (\mu, 0)$$

(3.6)

and the vacuum energy is $\mu^4 \sim \delta \Lambda^4$. The mass spectrum of the fluctuations is identified by expanding the superpotential around (3.6). At tree level, some fields get a mass of order $\mu$. The one loop Coleman-Weinberg potential [13] stabilizes all of the fields that remained massless except $Z^j_A$. Henceforth, it will be diagonalized using $SU(2)_L \times SU(2)_R$.\(^7\) At two loops, this pseudo modulus is destabilized [14–16]. For $Z^j_A \gg \mu$,

$$V^{(2)}_{\text{eff}} \sim -\frac{1}{(16\pi^2)^2} \mu^4 \left( \log \left| \frac{Z^1}{\mu} \right|^2 + \left( \log \left| \frac{Z^2}{\mu} \right|^2 \right)^2 \right)$$

(3.7)

is a good approximation to the effective potential.

As before, a term $\lambda \Lambda Z^1 L$ is added to $W'$ in attempt to stabilize the runaway, leading to

$$V_{\text{tree}} \supset (\lambda \Lambda)^2 |Z^1_A| + |\lambda \Lambda L^2 + \sigma^A \bar{\sigma}_1|^2 = (\lambda \Lambda)^2 \left( |Z^1_A|^2 + |L^2_A|^2 \right)$$

(3.8)

where the last equality holds since $\sigma, \bar{\sigma}$ are stabilized at the origin. As in the electric description, the tree level interaction does not lift all runaway directions unless $SU(2)_L$ is weakly gauged. In the limit $\Delta \gg 1$, the vacuum sits at the minimum of $V_{\text{tree}} + V^{(2)}_{\text{eff}}$ where

$$D \equiv \frac{1}{8} \left( 4|L_1 L_2|^2 + \left( |Z^1|^2 - |Z^2|^2 + |L_1|^2 - |L_2|^2 \right)^2 \right).$$

(3.9)

\(^6\)For the sake of clarity, the magnetic Yukawa coupling and other numeric factors of $O(1)$ are omitted.

\(^7\)As in the electric description, the $SU(2)_R$ flavor symmetry is broken by turning on $\lambda$. Future constraints on $|\lambda|$ ensure that this symmetry breaking is small.
The general solution to \( \mathcal{D} = 0 \) has two possible parameterizations:

\[
L = \left( 0, \sqrt{|Z_1|^2 - |Z_2|^2} \right), \quad L = \left( \sqrt{|Z_2|^2 - |Z_1|^2}, 0 \right).
\]  

(3.10)

Since there is no gauge symmetry left to connect the two solutions, \( L \) must be stabilized at the origin. Otherwise, the magnetic dual will fail to have the same global symmetry as the underlying theory. Indeed, this will turn out to be the case.

It is useful to absorb all constants by a redefinition of \( Z \):

\[
Z_i = \mu \sqrt{a} \zeta^{(i)}
\]

(3.11)

where \( a \equiv (16\pi^2)^{-2} 2\lambda^{-2} \delta^{1/2} \) and \( Z_A \gg \mu \) requires \( a \gg 1 \). Using these variables, the effective potential is

\[
V_{\text{eff}} = \frac{\mu^4}{(16\pi^2)^2} \left( 2 \zeta^{(1)}_\square + 2 \left| \zeta^{(2)}_\square - \zeta^{(1)}_\square \right| - \log \left[ a \zeta^{(1)}_\square \right]^2 - \log \left[ a \zeta^{(2)}_\square \right]^2 \right).
\]  

(3.12)

For a given \( \zeta^{(1)}_\square \), a minimum is found for \( \zeta^{(2)}_\square = \zeta^{(1)}_\square \equiv \zeta \). As promised, \( L \) is stabilized at the origin and one remains with

\[
V_{\text{eff}} \bigg|_{\zeta^{(1)}_\square = \zeta^{(2)}_\square} = \frac{\mu^4}{(16\pi^2)^2} \left( \zeta^2 - \log \left[ a \zeta^2 \right]^2 \right).
\]  

(3.13)

Expanding around large values of \( \zeta \), one finds

\[
\langle Z^i_A \rangle = \mu \sqrt{2 a \log |a|} \sim \delta^{\frac{1}{2}} \lambda^{-1} \Lambda \frac{1}{16\pi^2}.
\]  

(3.14)

For the magnetic description to be reliable, \( \langle Z^i_A \rangle \ll \Lambda \) or equivalently \( \delta \ll \lambda^2 \) must hold.

One might question whether this analysis is at all valid as the loop computations were not expanded around the true tree level vacuum since the “Yukawa” coupling \( \lambda \) was introduced after the fact. In this respect, self-consistency requires that the mass given to \( Z^1_A \) and \( L_A \) by turning on \( \lambda \) be much smaller than all other tree level masses: \( \lambda \Lambda \ll \mu \). Luckily, this condition is equivalent to \( a \gg 1 \).

The weak gauging of \( SU(2)_L \) poses another question: might the two loop effective potential itself be altered by a contribution from gauge fields? The reason this does not happen is that all fields charged under \( SU(2)_L \) have supersymmetric masses at tree level and so the gauge spectrum is supersymmetric at one loop and therefore cannot contribute to the effective potential at two loops.
The electric and magnetic descriptions are seen to be valid in non-overlapping regions of parameter space, $\lambda \ll \delta$ and $\delta \ll \lambda^2$ respectively. As advocated, this is a consistency check of the duality. Moreover, the fields which play a nontrivial role in the vacuum structure of the magnetic dual are closely related to the gauge invariants of the electric theory even though the source of the runaway differs significantly between the two descriptions.

4 Summary

• The vacuum structure of the 3-2 model was studied for generic values of the gauged $SU(2)_L$ coupling $g_2$ and the “Yukawa” coupling $\lambda$. A controllable vacuum is found in the regime $g_2 \lesssim \lambda$ even when the non-perturbative interactions are generated by an instanton.
• Two vector-like pairs of $SU(3)_C$ with mass $m = \sqrt{\delta}\Lambda$ were added to the theory, and the vacuum was found in the limit $g_2 \gg \lambda$ in two regions of parameter space. For $\lambda \ll \delta$ the vacuum is at large values of the microscopic fields whereas for $\lambda^2 \gg \delta$ these fields confine and the vacuum is described by the magnetic dual theory. In both regions, the combined efforts of a tree level superpotential and the weak gauging of $SU(2)_L$ stabilize a runaway potential. The fact that the two regions of validity do not overlap is a consistency check of the duality. In the case studied here, it depends nontrivially on non-holomorphic data.

For $\lambda^2 \lesssim \delta \lesssim \lambda$, the vacuum is nowhere to be found as both descriptions break down; what happens here remains an open question.

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