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HIERARCHY GENERATION IN COMPACTIFIED SUPERSYMMETRIC MODELS

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ABSTRACT

The problem of generating a large hierarchy in compactified supersymmetric models is re-examined. It is shown how, even for the class of models for which $\text{Str} M^2$ is non-vanishing, a combination of non-perturbative effects and radiative corrections may lead to an exponentially large hierarchy. A corollary is that the couplings of the effective field theory in the visible sector should be small, i.e., perturbation theory should be applicable.

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Although superstring theories offer the exciting prospect of a consistent, finite unification of the fundamental forces [1], their implications for the physical world remain uncertain. The reason is that there must be several stages of symmetry breaking to go from the multidimensional space in which the string is defined to the four-dimensional world of our everyday experience. There must be a stage of compactification at a scale, Λ_G , at which all but four space-time dimensions are compactified with dimension $1/\Lambda_G$. There must also be stages of gauge and supersymmetry breaking to reduce the symmetry to that of the standard model.

In this letter we re-examine the expectations for supersymmetry breaking in such models, paying particular regard to the question of generating the mass hierarchy. Although supersymmetry in the effective low-energy theory after compactification ensures the stability of a hierarchy through the non-renormalization theorems there remains the question of how the hierarchy is generated and what sets its scale. Attempts to answer this have utilized flat directions, in which the effective scalar potential is flat at the tree level in some field direction, and used radiative corrections to determine the previously unknown field scale [2-5]. Since these are driven by terms depending logarithmically on the scale of supersymmetry breaking, they may differ from other scales in the theory by exponentially small factors. In string theories, it is thought that the effective potential after compactification is of the "no-scale" type [6] which indeed possesses such flat directions at tree level.

Much work has gone into the determination of radiative corrections [2-5]. In principle, these should be calculated by evaluating the full string theoretic vacuum amplitude. This certainly includes a sum over all massive string excitations^{*)} but in practice we are not able to perform the full string-theoretic calculation in a realistic model with broken supersymmetry and so we start by ignoring all the massive states, both the string excitations and the Kaluza-Klein states which may be generated by compactification, retaining only those present in the effective low-energy four-dimensional field theory. Because of this truncation we should not despair when we find an apparently non-zero cosmological constant at the minimum of our one-loop effective potential. We will return to a discussion of the effects of the heavy modes later.

The calculation of the one-loop effective potential for the massless modes after compactification gives the general form [7]

^{*)} But note it is not equivalent to it, since the invariant domain of moduli space over which one must integrate is smaller than that to which one would be led by a naïve sum over states.

$$V_{\text{eff}}^{(1\text{-loop})} = \frac{1}{64\pi^2} \text{Str} \left\{ \Lambda^2 M^2 + \Lambda^4 \ln\left(1 + \frac{M^2}{\Lambda^2}\right) - M^4 \ln\left(1 + \frac{\Lambda^2}{M^2}\right) \right\} \quad (1)$$

where $\text{Str}\{ \}$ denotes the supertrace sum and Λ^2 denotes the maximum value of the virtual loop momentum squared for which the supersymmetry breaking effects are non-zero. Since the dominant contribution to V_{eff} comes from the upper value of this virtual loop momentum, it is important to determine the value of Λ^2 , and we will return to this question in the context of a specific model shortly. In general, however, one expects $\text{Str}(M^2) \ll \Lambda^2$ since, after taking the supertrace in Eq. (1), it is only the supersymmetry breaking contributions to M^2 that remain and we are looking for solutions to the hierarchy problem in which these terms are much smaller than the other scalar in this theory. For $\text{Str}(M^2) \approx m_{3/2}^2 \ll \Lambda^2$ we may expand Eq. (1) in the form

$$V_{\text{eff}}^{(1\text{-loop})} = \frac{1}{32\pi^2} \Lambda^2 \text{Str} \{ M^2 \} + \frac{1}{64\pi^2} \text{Str} \left\{ M^4 \ln \frac{M^2}{\Lambda^2} \right\} + \dots \quad (2)$$

In previous suggestions for using radiative corrections to fix the classically undetermined scale, either the leading (in the limit $M^2 \ll \Lambda^2$) term $\propto \Lambda^2 \text{Str} M^2$ has been assumed to vanish [6,8], or the full form (1) of $V_1(M, \Lambda)$ has been used [2-5]. There are indeed models where $\text{Str} M^2$ or even $\text{Str}(M^2)$ vanishes [9], but the generic situation is that the $\Lambda^2 \text{Str} M^2$ term dominates:

$$V_1(M, \Lambda) \approx \frac{\Lambda^2}{32\pi^2} \text{Str} \{ M^2 \} \quad (3)$$

The key idea of this paper is to use (3) to fix the classically undetermined scale.

At first sight, this approach looks unpromising because either $\text{Str} M^2 < 0$, in which case the potential is unstable in the region of $M^2 \ll \Lambda^2$ and one expects M^2 to be fixed at $O(\Lambda^2)$, or $\text{Str} M^2 > 0$, in which case there is no non-trivial minimum.

However once non-perturbative effects are included the situation changes, for these effects may lift the flat directions causing the potential to become unstable. In this case the terms of Eq. (3) may stabilize the potential giving rise to a non-trivial minimum for $\text{Str} M^2 > 0$ and it is possible for a large hierarchy to be generated with $\text{Str} M^2 \ll \Lambda^2$ as is phenomenologically desirable.

Let us turn now to a realization of this idea in the context of superstring models. After compactification the effective potential for a one-generation model is thought to have the "no-scale" form [10] [here we assume an $SU(N,1)$ generalization [6] for N generations]. To accommodate non-perturbative effects we allow for an S and a T dependence in the superpotential, where S and T are the usual complex chiral superfields related to the dilaton and the field setting the scale of compactification. The novel feature here is the inclusion of the T dependent piece. We have

$$V^{(0)} = \frac{1}{8S_R T_R'^3} \left\{ \left| W - 2S_R \frac{\partial W}{\partial S} \right|^2 + \frac{T_R'}{3} \left| \frac{\partial W}{\partial \phi} + 2\phi \frac{\partial W}{\partial T} \right|^2 + \frac{4T_R'}{3} \left| \frac{\partial W}{\partial T} - \frac{3W}{2T_R'} \right|^2 - 3|W|^2 \right\} \quad (4)$$

where

$$2T_R' \equiv T + \bar{T} - 2|\phi|^2 \quad ; \quad 2S_R \equiv S + \bar{S} \quad ; \quad 2T_R \equiv T + \bar{T} \quad (5)$$

following from the Kähler potential G given by

$$G = -\ln(2S_R) - 3\ln(2T_R') - \ln|W(S, T, \phi)|^2 \quad (6)$$

We assume that supersymmetry is broken in the hidden sector by non-perturbative effects giving rise to a gaugino condensate, $\langle \lambda\lambda \rangle$ [11]. Then the superpotential $W(S, T, \phi)$ will be of the form

$$\begin{aligned} W(S, T, \phi) &= P(T, \phi) + c + S \langle \lambda\lambda \rangle \\ &= P(T, \phi) + c + h \exp(-3S/2b_0) \end{aligned} \quad (7)$$

Here $P(T, \phi)$ is the superpotential describing the interaction of the chiral superfields, $h \exp(-3S/2b_0)$ arises as a result of gaugino condensation [11] and is given by the expectation value $\langle S\bar{\lambda}\lambda \rangle$ where we have used the form of the gauge kinetic function $f_{\alpha\beta} = \delta_{\alpha\beta} S$.

In Eq. (7) c is a contribution which may come from a non-vanishing VEV of the three index antisymmetric tensor field in the graviton supermultiplet [10] and/or

from a non-vanishing VEV of the superpotential [5]^{*}). In the case that it arises due to a VEV for the three index antisymmetric tensor it is expected to be quantized in Planck units $c \propto M_{\text{Planck}}$ [13]. For the purposes of generating a large hierarchy this proves to be unacceptably large, so we turn to the other possibility, namely that c is generated as the VEV of the superpotential [5] in which case it may be much smaller. For example, if $P(N) = aN^3 + MN^2$, where N is a singlet chiral superfield, there are two flat directions giving $F_N = 0$, namely $N = 0$ or $N = -2M/3a$. Taking the latter non-trivial solution⁺) gives $c = \langle P(N) \rangle = 4M^3/27a^2$. Now the magnitude of c is determined by the magnitude of M^3/a^2 . In the absence of non-perturbative effects, one expects massless singlet chiral fields after compactification, such as the moduli and the $H_1(\text{End}T)$ fields [14] in Calabi-Yau compactifications. The latter will, in general, acquire non-zero mass non-perturbatively [15] and for them $M \propto \Lambda_G \exp(-a'/(a'^2 \Lambda^2))$ where Λ_G is the compactification scale ($4\pi m_p^2/\sqrt{ST}$) and $(a')^{-1}$ is the string scale ($2m_p^2/\sqrt{S}$) [16], i.e.,

$$\frac{1}{a'^2 \Lambda_G^2} = \frac{1}{\pi} \cdot \frac{T}{m_p} \quad (8)$$

where $m_p^2 = 1/8\pi G_N$ will be our fundamental mass unit. The parameter a' has been calculated in certain orbifold and associated Calabi-Yau models where it is found that $0 < a' < 1$ [17].

Using this form for M shows that $c^{1/3}$ may be much smaller than the Planck scale provided only that Λ_G be somewhat less than M_p . However, the relationship of Eq. (8) shows that now M and hence the superpotential P will be S and T dependent so we need to use the full form of Eq. (4)^s) when determining the tree level potential including non-perturbative effects. In fact the other couplings, $eg\lambda$, in the superpotential may similarly be S and T dependent. At this stage it is useful to identify the free parameters involved in the effective potential. In Eqs. (4)

^{*}) Kounnas and Porrati [12] have argued against the presence of a contribution c , on the grounds that it violates an exact non-anomalous symmetry of the theory. However, as the symmetry is spontaneously broken by the VEVs of scalar fields, this term will occur as the VEV of the superpotential. Similarly we do not see why the VEV of the antisymmetric tensor should not spontaneously break the symmetry and generate a c term. We do not include here the possible effects of modification of the kinetic scalar term discussed in [9] as they do not significantly alter the analysis.

⁺) Which minimum is chosen depends on initial conditions which are determined by thermal corrections. In this simple example the minimum of the high temperature potential lies at $N = -M/3a$ so both minima are possible. More complicated examples may be constructed uniquely choosing the non-trivial minima.

^s) I thank P. Binétruy and I. Hinchliffe for emphasizing this point to me.

and (7) there are four parameters namely s, t, c and h . The quantity h is introduced in the parametrization of the gaugino condensate [cf. Eq. (7)]. Its exact value is determined by the strong interaction dynamics in the hidden sector, as discussed in detail at the end of the paper. For present purposes we just need the result that h is a constant^{*}) of $O(M_p^3)$. Two combinations of the remaining parameters are fixed

by minimizing the first two terms of Eq. (4) - for convenience we may take them to be s and c (or equivalently $\langle N \rangle$). The remaining parameter, T , may be fixed by minimizing the latter two terms of Eq. (4). If $-\text{Re}\{W^* \partial W / \partial T\} > 0$, these terms are positive and the minimum corresponds to $T \rightarrow \infty$ or $\Lambda_G \rightarrow 0$, clearly unacceptable. On the other hand, if $-\text{Re}\{W^* \partial W / \partial T\} < 0$, the last two terms of Eqs. (5) are negative at the minimum and it will correspond to $T \rightarrow 0$ or $\Lambda_G \rightarrow \infty$, again unacceptable. However, in this latter case radiative corrections may stabilize the potential and determine T and hence Λ_G at a finite value.

To demonstrate this latter possibility we work with a superpotential of the form

$$\begin{aligned} \mathcal{P}(T, N) &= \lambda(T) N^3 + M(T) N^2 \\ &\equiv \lambda_0(T) \exp(-\beta T/m_T) N^3 + M_0(T) \exp(-\alpha T/m_P) N^2 \end{aligned} \quad (9)$$

where we have explicitly exhibited the leading exponential behaviour of M (a, b are positive) and we have also allowed for the possibility that the leading contribution to λ is also non-perturbative in origin. In addition we have allowed for a residual T dependence in $M_0(T)$ and $\lambda_0(T)$. This is the form found by explicit calculation in a Calabi-Yau model close to the orbifold limit [17]. There the T dependence of M_0 and λ_0 arises because they are given by a power series in $\langle b_i \rangle / (\Lambda_G^2 \alpha'^2)$ where the b_i are the blowing up modes or moduli of the Calabi-Yau model.

Now we may compute $\text{Re}\{W^* \partial W / \partial T\}$ subject to the condition that the second term in Eq. (4) vanishes. Although there are several possible solutions of the latter equation, we will be interested only in the solution close to the global

^{*}) Binétruy et al. [4] impose the condition $\partial V / \partial h = 0$. We note that this condition follows automatically from the minimization of dynamical field variables c, s, t if h is multiplicatively renormalized and that in this case, it is not necessary to include this condition explicitly.

metry minimum $\partial W/\partial \phi \approx 0$, for in this case we can arrive at a stable hierarchy with no fields acquiring VEVs much larger than the Planck scale. In this case we have $P \sim 1/3 M N^2 \approx (4/27)(M^3/\lambda^2)$ at minimum and we find

$$-Re\left\{W^* \frac{\partial W}{\partial T}\right\} = -3|W|^2 \left(1 + \frac{2b_0}{3S}\right) \left\{ \left(\frac{2}{3}\beta - \alpha\right) \frac{1}{m_P} - \left(\frac{2}{3} \frac{\lambda'_0(T)}{\lambda_0(T)} - \frac{M'_0(T)}{M_0(T)}\right) \right\} \quad (10)$$

Since the gravitino mass is given by

$$m_{3/2} = \frac{1}{8S_R T_R^{1/3}} |<W>|^2 \quad (11)$$

we see that the condition that $m_{3/2}$ should be suppressed by an exponentially small factor is that $(\frac{2}{3}\beta - \alpha)$ be negative [here we have used $W \propto P \propto M^3/\lambda^2$ and the form of Eq. (9) for M and λ]. In this case the first term of Eq. (10) is positive but the overall sign may be negative (corresponding to a negative potential at tree level) due to the second set of terms in Eq. (10). In fact, we expect this to happen for, as noted above, they are determined by some function of $(\langle b_i \rangle T)$ and the VEVs of the moduli, b_i , will adjust, if possible, to give the minimum energy solution. Following the analysis of Ref. [18] we note that these moduli will also appear in the potential via the second term of Eq. (4) with $\phi = b_i$. If we take just one such field to illustrate, we find from Eqs. (4), (10) and the result that λ and M depend on bT :

$$V^{(0)} \approx \frac{9|W|^2 \left(1 + \frac{2b_0}{3S}\right)^2 y^2}{8S_R T_R'^2} \left\{ 12|b|^2 \left|1 - \frac{yT}{4b^2}\right|^2 + \frac{4}{3}T - \frac{4}{3y} \left(1 + \frac{2b_0}{3S}\right)^{-1} \right\} \quad (12)$$

where

$$y = \left(\frac{2}{3}\beta - \alpha\right) \frac{1}{m_P} - \left(\frac{2}{3} \frac{\lambda'_0(Tb)}{\lambda_0(Tb)} - \frac{M'_0(Tb)}{M_0(Tb)}\right) b \quad (13)$$

and we have already used the variation with respect to S to set the first term of Eq. (4) to zero. Clearly, Eq. (12) becomes negative for small $y(b)$ as anticipated, fixing the VEV of b . Then, as discussed above, minimization with respect to T gives $T \rightarrow 0$ or $\Lambda_G \rightarrow \infty$, and we must look at radiative corrections to find if the potential is stabilized.

Thus we turn now the determination of the one-loop effective potential, Eq. (1). As discussed above, the dominant contribution comes from the region of maximum virtual loop momentum. For momentum scales above the compactification scale, however, we expect the spectrum of the theory to be that of the uncompactified theory, for the forces giving rise to the compactified fields in the low-

energy effective theory become weak and the fields dissolve leaving the spectrum of the underlying theory. Since supersymmetry breaking is triggered by VEVs in compactified fields, we expect that these effects, too, will disappear for momentum scales above Λ_G . It follows therefore that the dominant contribution to Eq. (1) comes from virtual momentum scales less than but of order Λ_G and that Λ^2 in Eq. (1) should be identified with Λ_G^2 .

The light mass spectrum in this momentum range is easy to determine using the effective $N = 1$ supersymmetric Lagrangian involving the compactified zero modes. In the hidden sector the gauge bosons and gauginos are massless for we are at virtual momentum scales ($k \approx \Lambda_G > \Lambda_c$) above the condensation scale and so the gaugino condensate will disappear leaving unbroken the hidden sector gauge group. The gravitino has a mass given by Eq. (11) and contributes to the supertrace in Eq. (11).

Finally in the visible sector the scalar component, ϕ , of the gauge non-singlet chiral superfields all acquire a supersymmetry breaking contribution to their mass. Consider first the explicit ϕ dependence coming from the second term of Eq. (5). Allowing for the non-canonical kinetic energy, we find the contribution

$$m_\phi^{(0)2} = \frac{1}{6S_R T_R} \left| \frac{\partial W}{\partial \tau} \right|^2 = \frac{9g^2 \left(1 + \frac{2b_0}{3s}\right)^2 |W|^2}{6S_R T_R} \quad (14)$$

There is another contribution to the ϕ mass-squared through the $T'_R(\phi)$ dependence of Eq. (12) where $T'_R(\phi)$ is given in Eq. (5). This is negative and proportional to the term in curly brackets in Eq. (12). This we will ignore for we are interested in determining the maximum possible hierarchy and we will see any negative contribution to m_ϕ^2 reduces the hierarchy. Also we will argue presently that there are likely to be additional positive contributions to Eq. (12) coming from the sector driving the gaugino condensate and the net effect will be to make the term in curly brackets in Eq. (12) very small, corresponding to a small cosmological constant.

Note that the appearance of a scalar mass as given in Eq. (14) is in marked contrast to the case normally considered in which no chiral superfields acquire large VEVs for then the expectation value of the right-hand side of Eq. (14) vanishes, as it would if we have ignored non-perturbative effects.

The one-loop contribution, Eq. (4), is then

$$V'(M, \Lambda_G) = + \frac{\Lambda_G^2}{16\pi^2} \left(-m_{3/2}^2 + N m_\phi^{(0)2} \right) \quad (15)$$

where N is the number of scalar components of chiral superfields. Although the combination $V^{(0)} + V^{(1)}$ is a complicated function of T (or equivalently Λ_G) we may determine some general properties of the minimum from the general structure we have derived. The last term of Eq. (15) stabilizes the potential along the direction, $\partial W / \partial R = 0$, in which at tree level the potential was unbounded from below. It is straightforward to verify that, in general, there are no other unbounded minima for $V^{(0)} + V^{(1)}$; for example the direction $W = 0$, $\partial W / \partial T = 0$ would be favoured by the first term of Eq. (15) but is possible only through a choice of $\langle b \rangle$ which gives a positive contribution at tree level via the second term of Eq. (5).

Thus we conclude that $V^{(0)} + V^{(1)}$ possesses a non-trivial minimum. We can go further if this minimum corresponds to large S and T . Then keeping the dominant terms in Eq. (12) and minimizing with respect to b gives

$$V^{(0)} + V^{(1)} = 3 m_{3/2}^2(T) \left\{ -1 + \frac{N \Lambda_G^2(T)}{8\pi^2} \right\} \quad (16)$$

The minimum of this corresponds closely to the point at which it changes sign giving

$$\Lambda_G^2 \approx \frac{16\pi^2 m_P^2}{ST} \approx \frac{8\pi^2}{N} \quad (17)$$

Does this imply an hierarchy? To see this, note that for large S , T the minimization of the first term of Eq. (4) requires

$$\frac{3S}{2b_0} \approx (-2\beta + 3\alpha)T \quad (18)$$

Equations (17) and (18) imply that S and T are indeed large if $3/2b_0$ and $(2\beta - 3\alpha)$ are comparable in magnitude and N is large (as it is in viable models, justifying, a posteriori, the approximation used in obtaining Eq. (16)). The value of the gravitino mass is

$$m_{3/2}^2 = \frac{9}{8} \frac{\delta^2}{N} h^2 \exp(-3S/b_0) \quad (19)$$

where

$$\delta = \frac{2}{3} (-2\beta + 3\alpha) \quad ; \quad h \approx m_P$$

$$S = \sqrt{2N} \sqrt{\delta b_0} \quad (20)$$

For large N a large hierarchy may develop. For example, in the three generation Calabi-Yau manifold, it is found [19] that $N = 15 \times 27$ corresponding to $m_{3/2}^2 \sim m_p^2 e^{-90/(\delta/b_0)}$. Here b_0 is the hidden sector β -function coefficient; if the hidden sector group is E_8 , $b_0 = +90$ and we would not expect a large hierarchy for reasonable δ . On the other hand, the hidden sector group may be broken on compactification. If it is broken to $SU(3)$, then $b_0 = 9$ and a very large hierarchy may develop.

At the minimum of the potential, $V^{(0)} + V^{(1)}$, the energy is negative [cf. Eq. (16)] corresponding to a negative cosmological constant. In view of our neglect of Kaluza-Klein and massive string modes we should perhaps not worry about this, but there is another potential contribution that has so far been ignored. This is the contribution from hidden sector non-perturbative effects which must be responsible for the gaugino condensation. The last term of Eq. (7) partially parametrizes these effects but it would clearly be desirable to calculate these effects explicitly. A partial approach to this has been given in Refs. [20] and [21], where an effective Lagrangian describing the gaugino condensation was constructed. In this, one introduces a composite multiplet $U = WW$ which contains the gaugino composite $\lambda\lambda$ as its scalar component. The $N = 1$ effective supergravity Lagrangian is then given by adding to the Kähler potential, Eq. (6), the term $\gamma'(u^*u)^{1/3}$ and substituting for the last term of Eq. (7) the piece

$$\frac{1}{4} u s + \left(\frac{b_0}{6}\right) u \log(m_p^{-3} u)$$

The net effect is to add to the potential Eq. (5) the term

$$V^{(0)'} = \frac{1}{8 S_R T_R'^3} \left[\frac{9}{8} (u^* u)^{2/3} + \frac{1}{4} s + \frac{b_0}{6} (1 + \log(m_p^{-3} u)) + \frac{\gamma'}{3u} (u^* u) \right]^2 \quad (21)$$

Minimization of the full potential now includes minimization with respect to the gaugino condensate. The net result is that the gaugino condensate is fixed at $\sim m_p^3 \exp(-3S/2b_0)$ as before but there is an additional positive contribution to $V^{(0)}$ coming from the new terms in the superpotential and corresponding to a contribution to the potential energy from strong non-perturbative hidden sector effects. As a result the total potential energy at minimum may be zero as desired, while the hierarchy generation proceeds essentially as discussed above.

To summarize, we have found that, in the large N limit, the effective potential has a stable minimum at which, for large N , an exponentially small mass hierarchy may develop for the gravitino mass and other supersymmetry breaking masses. As a corollary to this, note that large S [cf. Eq. (20)] corresponds to small gauge coupling ($g_{GUT}^2 = m_p^2/S_R$). Moreover, after further intermediate stages of gauge symmetry breaking, we expect the 27 and $\overline{27}$ multiplets to pair up [22] so

that the effective theory at low energies will not have a large number of chiral superfields and hence the perturbative expansion will not contain large N factors. As a result we may expect the perturbative expansion in both the gauge and Yukawa couplings to be good in the effective low-energy theory.

Of course we should worry about terms not included in our analysis. We certainly expect corrections due to massive Kaluza-Klein and string modes. For large masses M , these terms are suppressed in their contribution to the one-loop potential by the factor (Λ_G^2/M^2) relative to light modes. However, the first level of Kaluza-Klein modes have $M^2 \approx \Lambda_G^2$, so their contribution is not negligible and we expect additional contributions to Eq. (15). Provided they do not change sign of this term, the mechanism for hierarchy generation remains. This just illustrates the point that the appearance of a large hierarchy may be expected in a wide class of supersymmetric models - indeed changing the form of the original Lagrangian in the manner discussed in Ref. [24] does not alter the structure which generates the hierarchy.

Finally, we consider the mass spectrum to be expected at low energies. As discussed above, scalars acquire a supersymmetry breaking mass at tree level $\propto m_{3/2}$ quite different structure from previously discussed supersymmetry breaking schemes. By contrast, gaugino masses are generated first in one loop order [23] and are expected to be smaller. Thus the spectrum to be expected at low energies may be much different from our present expectations. For any particular model, this spectrum may be obtained using the renormalization group equations in the usual way with the initial conditions those given above. However, the analysis will be complicated in large N theories, by the large number of Yukawa couplings so it is not possible to make general statements about the resulting spectrum.

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