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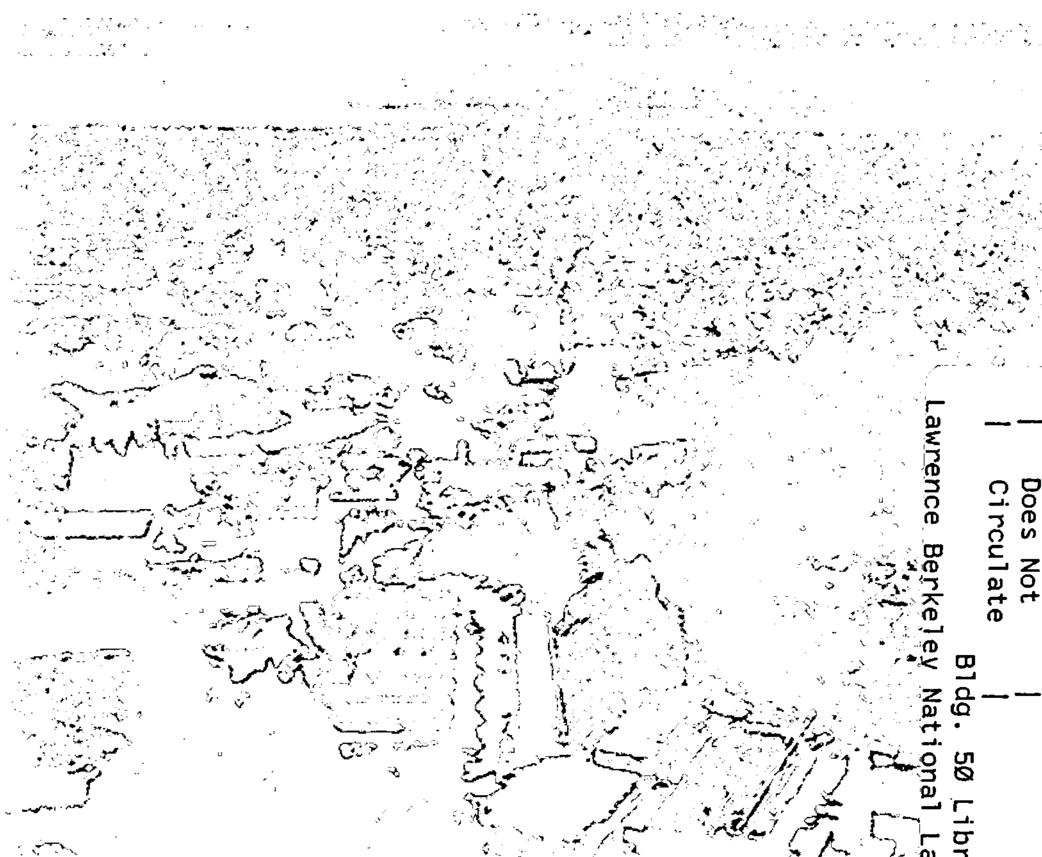

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Propagating Higgs Boundstates from Sfermions

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Abstract

A model of supersymmetric dynamical electroweak breaking with propagating sfermionic Higgs boundstates is constructed. The low energy effective theory is represented by a slight extension of the MSSM, including 2 additional Higgs doublets and neutrino Yukawa couplings. A large $tan\beta$ is a necessary condition. The model could be relevant in approaches which derive propagating Higgs boundstates from strings.

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In the framework of the minimal supersymmetric standard model (MSSM) the renormalization group running of the Yukawa couplings leads to Landau poles around the Planck scale for small or very large $tan\beta$. One possible way to interpret this fact is to understand the Higgs fields as propagating boundstates [1, 2, 3] that break up at the Planck scale. However, the separation of binding scale and electroweak scale is a major problem in a supersymmetric model of dynamical electroweak breaking. Due to the cancellation of quadratic contributions in SUSY it is not possible to produce a critical effective coupling from an underlying gauge interaction. It is possible to get the appropriate operators from supergravity [4], however those operators are highly nonlinear and do not look too attractive. The most promising framework to produce the wanted enhanced operators seems to be a strongly coupled stringy scenario with some large compactified dimensions [5]. There exist two alternatives to construct dynamical electroweak breaking in a supersymmetric framework. The possibility that has been considered so far [3] is dynamical breaking induced by a enhanced nonrenormalizable D–term interaction. This is a top–condensation scenario where the two effective Higgs doublets of the MSSM–like low energy theory are represented by a two–sfermion boundstate and a more complicated boundstate dominated by a top(bottom)–antitop contribution. General arguments in a stringy scenario however favour the second alternative, enhanced F–terms [5]. Therefore the goal of this letter is to formulate a model of F–term induced Higgs boundstates along the lines followed in the D–term induced case.

The basic idea of an effective model of dynamical electroweak breaking is to avoid fundamental Higgs fields and instead introduce a nonrenormalizable coupling term. If this coupling term is stronger than a critical coupling value, it induces electroweak breaking, effectively described by a Nambu–Jona-Lasinio type gap equation [6]. A finetuned model of electroweak breaking separates the binding scale of the new interaction term, represented by a high energy cutoff Λ in the effective theory, from the electroweak scale. The Nambu–Jona-Lasinio description is an excellent approximation in this case. In a supersymmetric model the separation of scales is realized by a scale difference between Λ and the soft breaking scale Δ plus an enhancement of the nonrenormalizable coupling G from its natural value $1/\Lambda^2$ to $1/\Delta^2$ [3].

An important role in finetuned dynamical electroweak breaking is played by the renormalization group approach [1]. This approach uses the fact that a model of finetuned dynamical electroweak breaking is well described up to the cutoff scale by a low energy effective theory with an electroweak breaking scalar sector (e.g. the SM, MSSM, etc.). Therefore it is possible to define the low energy characteristics of dynamical breaking by identifying the low energy Lagrangian with the fundamental Lagrangian of the model at the cutoff scale. This is most easily done by writing the nonrenormalizable interaction of the fundamental Lagrangian in an auxiliary field formalism. The auxiliary fields – modulo some normalization factor – can then be directly identified with the scalars of the effective theory. This identification gives certain conditions for the renormalization group running of the parameters of the low energy effective theory (the so called constituent conditions). As it is argued in detail in [1, 7], provided there is no new physics below the cutoff scale an electroweak breaking low energy theory that fulfills these conditions necessarily implies dynamical symmetry breaking induced by new physics at the cutoff scale. Our discussion will take place entirely in the framework of the renormalization group approach. We will write down the auxiliary field Lagrangian of F–term induced dynamical electroweak breaking, show, that the constituent conditions are fulfilled and extract the basic properties of the model.

For reasons of comparison we start with D–term induced electroweak breaking (SUSY top–condensation). There the Lagrangian has the form:

$$\begin{aligned}\mathcal{L}_D = \mathcal{L}_{YM} \quad &+ \quad \int d^2\theta d^2\bar{\theta}(\bar{Q}e^{2V_Q}Q + T^c e^{-2V_T}\bar{T}^c + B^c e^{-2V_B}\bar{B}^c)(1 - \Delta^2\theta^2\bar{\theta}^2) \\ &+ \quad G_T \int d^2\theta d^2\bar{\theta}[(\bar{Q}\ \bar{T}^c)e^{2V_Q - 2V_T}(QT^c)](1 - 2\Delta^2\theta^2\bar{\theta}^2 + \delta\bar{\theta}^2 + \delta\theta^2)\ ,\end{aligned} \tag{1}$$

where $\mathcal{L}_{YM}$ contains the usual SUSY kinetic terms for gauge fields, Q (T^c, B^c) are SU(2) doublet (singlet) chiral quark superfields. Δ^2 and δ are SUSY soft breaking parameters. Throughout this work superfields will be denoted by capital letters and component fields by small letters except for the vectorfield which is identifiable by its Lorentz index. Reformulated in auxiliary fields eq.(1) corresponds to:

$$\begin{aligned}\mathcal{L}_D = \mathcal{L}_{YM} \quad &+ \quad \int d^2\theta d^2\bar{\theta}(\bar{Q}e^{2V_Q}Q + T^c e^{-2V_T}\bar{T}^c + B^c e^{-2V_B}\bar{B}^c)(1 - \Delta^2\theta^2\bar{\theta}^2) \\ &+ \quad \int d^2\theta d^2\bar{\theta}\ \bar{\hat{H}}_1 e^{2V_{H_1}}\hat{H}_1(1 - M_{\hat{H}}^2\theta^2\bar{\theta}^2) \\ &- \quad \int d^2\theta\epsilon_{ij}(\mu_0\hat{H}_1^i\hat{H}_2^j(1 + \hat{B}\theta^2) - \hat{g}_T\hat{H}_2^j Q^i T^c) \\ &- \quad \int d^2\bar{\theta}\epsilon_{ij}(\mu_0\bar{\hat{H}}_1^i\bar{\hat{H}}_2^j(1 + \hat{B}\bar{\theta}^2) - \hat{g}_T\bar{T}^c\bar{Q}^i\bar{\hat{H}}_2^j)\ ,\end{aligned} \tag{2}$$

with $M_{\hat{H}}^2 = 2\Delta^2 + \delta^2$, $\hat{B} = -\delta$, $V_{H_1} = V_Q - V_T$ and $G_T = \frac{\hat{g}_T^2}{\hat{\mu}^2}$. The hat will always denote auxiliary fields respectively their couplings. As already mentioned the auxiliary fields in eq.(2) represent the two MSSM Higgs superfields at the cutoff scale Λ in a different normalization. Note that not all kinetic terms of the composite scalar fields vanish at the cutoff scale. The auxiliary field $\hat{H}_1$ has a kinetic term. This kinetic term is unavoidable in an auxiliary field formulation of a supersymmetric four–fermion theory for the following reason: Describing a supersymmetric four–fermion interaction requires supersymmetric auxiliary field terms. These have to include fermion couplings to the auxiliary scalar as well as an auxiliary scalar mass term. Scalar mass terms are gained from the μ–term $\int d^2\theta \epsilon_{ij} \mu_0 \hat{H}_1^i \hat{H}_2^j$ by integrating out $\hat{F}_H$–fields. For integrating out one needs $\hat{F}_H^2$–terms which come from the same superfield as the scalar kinetic term. Thus a scalar kinetic term is necessary for implementing the auxiliary field formalism for a four fermion interaction. Of course the second auxiliary kinetic term must not occur in order to keep the auxiliary character of the H–fields. In the dynamical picture this relates to the fact that the second scalar is not a fermion-boundstate but a two–scalar boundstate and is produced directly by using the auxiliary field $\hat{F}_{H_2}$ as a Lagrangian multiplier.

To get a better understanding of the structure of the auxiliary field concept, it is instructive to have a more general look at the possibilities to construct an auxiliary field Lagrangian: We start with an MSSM–like Lagrangian (forgetting about gauge fields and soft breaking terms at the moment) which involves two Higgses, both possessing kinetic terms and Yukawa couplings.

$$\begin{aligned}\mathcal{L} &= \int d^2\theta d^2\bar{\theta}\, (\bar{H}_1 H_1 + \bar{H}_2 H_2) - \int d^2\theta (\epsilon_{ij}(\mu H_1^i H_2^j - g_T H_2^j Q^i T^c) - g_B H_1^j Q^i B^c) \\ &- \int d^2\bar{\theta} (\epsilon_{ij}(\mu \bar{H}_1^i \bar{H}_2^j - g_T \bar{T}^c \bar{Q}^i \bar{H}_2^j) - g_B \bar{B}^c \bar{Q}^i \bar{H}_1^j)\ . \end{aligned} \tag{3}$$

Now we want to integrate out the massive fields H_1 and H_2. To do this in superfield formalism we write the D–terms in eq.(3) as F–terms

$$\begin{aligned}\int d^2\theta d^2\bar{\theta}\ \bar{H}_1 H_1 &= \int d^2\theta \bar{D}_\alpha^2 \bar{H}_1 H_1 \\ \int d^2\theta d^2\bar{\theta}\ \bar{H}_2 H_2 &= \int d^2\theta \bar{D}_\alpha^2 \bar{H}_2 H_2 \end{aligned} \tag{4}$$

and are able to integrate out the H–fields under the integral $\int d^2\theta$. The Euler–Lagrange equations are

$$H_1 = \frac{g_T}{\mu}(QT^c - \frac{1}{4}\bar{D}_\alpha^2 \bar{H}_2) \tag{5}$$

$$H_2 = \frac{g_B}{\mu}(QB^c - \frac{1}{4}\bar{D}_\alpha^2 \bar{H}_1) \tag{6}$$

Re-insertion of H_1 and H_2 eliminates these fields up to order $\frac{1}{\mu^2}$. The resulting effective Lagrangian has the form

$$\begin{aligned}\mathcal{L}_{eff} &= \int d^2\theta d^2\bar{\theta}[G_T(\bar{Q}\bar{T}^c)(QT^c) + G_B(\bar{Q}\bar{B}^c)(QB^c)] \\ &+ G\int d^2\theta QB^cQT^c + G\int d^2\bar{\theta}\ \bar{Q}\bar{B}^c\bar{Q}\bar{T}^c + \text{higher orders}\ ,\end{aligned} \tag{7}$$

where we have rewritten the SUSY derivatives D_α as θ–integrals again and used $G = \frac{g_T g_B}{\mu}$, $G_T = \frac{g_T^2}{\mu}$ and $G_B = \frac{g_B^2}{\mu}$. We get F–terms stemming from the μ–term and D–terms stemming from the kinetic terms of $\hat{H}_{1(2)}$ The propagation of the heavy H–fields is secluded in the SUSY derivatives of the higher order contributions. Now we remove one kinetic term in eq.(3) and get a Lagrangian

$$\begin{aligned}\mathcal{L}' &= \int d^2\theta d^2\bar{\theta}\ \bar{\hat{H}}_1\hat{H}_1 - \int d^2\theta\epsilon_{ij}(\hat{\mu}\hat{H}_1^i\hat{H}_2^j - \hat{g}_T\hat{H}_2^jQ^iT^c) - \hat{g}_B\hat{H}_1^jQ^iB^c) \\ &- \int d^2\bar{\theta}\epsilon_{ij}(\hat{\mu}\bar{\hat{H}}_1^i\bar{\hat{H}}_2^j - \hat{g}_T\bar{T}^c\bar{Q}^i\bar{\hat{H}}_2^j) - \hat{g}_B\bar{B}^c\bar{Q}^i\bar{\hat{H}}_1^j)\ .\end{aligned} \tag{8}$$

Eq. (5) reduces to

$$\hat{H}_1 = \frac{\hat{g}_T}{\hat{\mu}}QT^c\ . \tag{9}$$

Therefore there cannot exist any effective operators of dimension higher than 6, the contributions of propagation for *both* Higgs superfields vanish. We are confronted with an auxiliary field Lagrangian (which is why the H fields suddenly got hats), the kinetic term of $\hat{H}_1$ is just responsible for the derivative couplings in the D–term. The resulting interaction Lagrangian is

$$\mathcal{L}' = G_T\int d^2\theta d^2\bar{\theta}[(\bar{Q}\bar{T}^c)(QT^c)] + G\int d^2\theta QB^cQT^c + G\int d^2\bar{\theta}\bar{Q}\bar{B}^c\bar{Q}\bar{T}^c\ . \tag{10}$$

Now we can set to zero the second kinetic term in eq. (8) which leaves us with a purely holomorphic Lagrangian

$$\begin{aligned}\mathcal{L}'' = &- \int d^2\theta\epsilon_{ij}(\hat{\mu}\hat{H}_1^i\hat{H}_2^j - \hat{g}_T\hat{H}_2^jQ^iT^c) - \hat{g}_B\hat{H}_1^jQ^iB^c) \\ &- \int d^2\bar{\theta}\epsilon_{ij}(\hat{\mu}\bar{\hat{H}}_1^i\bar{\hat{H}}_2^j - \hat{g}_T\bar{T}^c\bar{Q}^i\bar{\hat{H}}_2^j) - \hat{g}_B\bar{B}^c\bar{Q}^i\bar{\hat{H}}_1^j)\end{aligned} \tag{11}$$

that of course corresponds to a holorphic interaction Lagrangian, i.e. a Lagrangian without any D–term:

$$\mathcal{L}'' = G\int d^2\theta QB^cQT^c + G\int d^2\bar{\theta}\bar{Q}\bar{B}^c\bar{Q}\bar{T}^c \tag{12}$$

The other possibility is to forbid one Yukawa term which corresponds to forbidding the four-superfield F–terms and brings us back to the SUSY top condensation Lagrangian of the type eq. (2). If we forbid both, kinetic terms for Hs and one Yukawa term we get zero for obvious reasons.

The Lagrangian of eq.(12) plus appropriate soft breaking terms is exactly the interaction Lagrangian we are looking for. The next step is to check whether this Lagrangian can be identified with the MSSM Lagrangian at some scale Λ by applying the constituent conditions. One can translate the vanishing of kinetic terms in the auxiliary Lagrangian into the low energy effective theory (the MSSM modulo possible extensions) by defining $H = g_Y\hat{H}$ and demanding a pole for the Yukawa coupling g_Y in the MSSM Lagrangian.

At this point we have to look more carefully at the role of the Landau pole in this framework: Of course, if we admit the small Yukawa couplings in our effective picture, eventually the top Yukawa pole will drive all Yukawa couplings into the same pole. However this effect is due to loop contributions which involve the Higgs doublet. Since the Higgs is a boundstate, in the full theory these contributions are a secondary effect compared to pure fermionic loop contributions. If one wants to have any control over the strong dynamics of electroweak breaking, one has to require that these secondary effects are suppressed in the strong coupling regime of the full theory and therefore irrelevant for the structure of dynamical breaking. This requirement is equivalent to demanding the validity of a $1/N_{c'}$ expansion. (See the comments on $1/N_{c'}$ later on.) In the large $N_{c'}$ limit the Higgs–loop contributions vanish. We can therefore distinguish between poles which still exist in the large $N_{c'}$ limit and poles which do not. Only the first type constitutes a constituent condition in a model of dynamical symmetry breaking. And we will always refer to the first type when mentioning a Landau pole henceforth.

In the case of SUSY top–condensation, where only one kinetic term vanished, there was just one Landau pole for g_T. Now, if we want to identify the low energy effective Lagrangian with the Lagrangian in eq.(11) at the scale Λ, we need two Landau poles for the Yukawa couplings g_B and g_T respectively. This however requires similar values for the two Yukawa couplings at the electroweak scale, which again implies a large $tan\beta$ to produce the correct mass ratio

m_t/m_b. Large values for $tan\beta$ lead to a Landau pole for the Yukawa couplings around the Planck scale since the contributions of a large g_B additionally increase the coupling values [8]. Thus a large $tan\beta$ scenario is compatible with a constituent Higgs picture. A more detailed picture of the running behaviour of the Yukawas depends on the SUSY breaking structure and the embedding into a GUT which to consider is beyond the scope of this letter.

An additional complication arises concerning the leptonic sector. In top condensation models usually the top is the only SM fermion that gets its mass directly from the dynamical breaking procedure, all other fermion masses, for quarks as well as for leptons, are fed down from the top mass in some not further specified pattern of higher order corrections. This seems to be plausible to some extent since the other fermion masses are altogether considerably suppressed against the top mass. In our new scenario however, the bottom quark gets its mass directly from dynamical breaking as well, the mass difference between top and bottom is now caused by a specific form of the vacuum and not by a different status in the process of symmetry breaking. This is in principle the more natural approach since it does not depend on treating differently two fields of the same generation. However it becomes highly implausible now to treat the τ–mass as a higher order effect compared to the bottom mass since both not only belong to the same generation but also have approximately the same value. Therefore it seems to be most natural to introduce an additional leptonic nonrenormalizable F–term to get the τ–mass directly from dynamical breaking as well. To do so, we have to couple the τ–superfield to some partner to get an F–term that preserves hypercharge. One could couple it to the top supermultiplet again but the much more elegant choice is to couple it to the τ–neutrino avoiding unnecessary asymmetries in the structure of the theory. The missing of an effective neutrino Yukawa coupling seems even more unnatural in the light of the recent experimental neutrino mass signatures. If the neutrinos are massive, there should exist low energy neutrino Yukawa couplings which produce Dirac masses at the electroweak scale plus high scale right handed Majorana masses leading to suppressed neutrino mass eigenstates via the a see–saw mechanism. The τ–neutrino Yukawa coupling naturally should have a value similar to the other third generation Yukawa couplings and therefore in our scenario produce another independent Landau pole respectively a fourth Higgs boundstate. In the large $tan\beta$ region a large neutrino Yukawa coupling does not have a big influence on the Yukawa unification properties [9]. Thus we define the interaction structure of our theory by the auxiliary Lagrangian

$$\mathcal{L}_F \;=\; \mathcal{L}_{YM} + \int d^2\theta d^2\bar{\theta}(\bar{Q}e^{2V_Q}Q + T^c e^{-2V_T}\bar{T}^c + B^c e^{-2V_B}\bar{B}^c)(1 - \Delta^2\theta^2\bar{\theta}^2)$$

$$
\begin{aligned}
&- \int d^2\theta\epsilon_{ij}(\hat{\mu}\hat{H}_1^i\hat{H}_2^j(1+\hat{B}\theta^2) - \hat{g}_T\hat{H}_2^jQ^iT^c - \hat{g}_B\hat{H}_1^jQ^iB^c) \\
&- \int d^2\bar{\theta}\epsilon_{ij}(\hat{\mu}\bar{\hat{H}}_1^i\bar{\hat{H}}_2^j(1+\hat{B}\bar{\theta}^2) - \hat{g}_T\bar{T}^c\bar{Q}^i\bar{\hat{H}}_2^j - \hat{g}_B\bar{B}^c\bar{Q}^i\bar{\hat{H}}_2^j) \\
&- \int d^2\theta\epsilon_{ij}(\hat{\mu}'\hat{H}_3^i\hat{H}_4^j(1+\hat{B}'\theta^2) - \hat{g}_N\hat{H}_4^jL^iN^c - \hat{g}_E\hat{H}_3^jL^iE^c) \\
&- \int d^2\bar{\theta}\epsilon_{ij}(\hat{\mu}'\bar{\hat{H}}_3^i\bar{\hat{H}}_4^j(1+\hat{B}'\bar{\theta}^2) - \hat{g}_N\bar{N}^c\bar{L}^i\bar{\hat{H}}_4^j - \hat{g}_E\bar{E}^c\bar{L}^i\bar{\hat{H}}_3^j)
\end{aligned}
\tag{13}
$$

The most general interaction Lagrangian would include also Yukawa terms of the type $\hat{g}'_E\hat{H}_1^jL^iE^c, \hat{g}'_B\hat{H}_3^jQ^iB^c, \hat{g}'_T\hat{H}_4^jQ^iT^c, \hat{g}'_N\hat{H}_2^jL^iN^c$ (or, equivalently, the μ–terms $\hat{\mu}''\bar{\hat{H}}_1^i\bar{\hat{H}}_4^j$ and $\hat{\mu}'''\bar{\hat{H}}_2^i\bar{\hat{H}}_3^j$). These terms would correspond to mixed lepton–quark four–superfield couplings. All arguments of the following discussion apply also in this more general case. We will stick to the simplest case for the sake of transparency. The introduction of trilinear soft breaking terms in eq.(13) would just correspond to a different parameterization of the same four–superfield interaction. The utility of the chosen parameterization will turn out later. To identify the Lagrangian eq.(13) with a – slightly extended – MSSM Lagrangian at the cutoff Λ we use four separate pole conditions for MSSM parameters.

$$
g_T^{-1}(\mu),\ g_B^{-1}(\mu),\ g_N^{-1}(\mu),\ g_E^{-1}(\mu)\ \rightarrow 0 \quad \text{for} \quad \mu \rightarrow \Lambda \tag{14}
$$

However to really identify the two Lagrangians it is necessary that the terms $\frac{\mu_1}{g_Tg_B}H_1^iH_2^j$ and $\frac{\mu_2}{g_Eg_N}H_3^iH_4^j$ do not vanish at the cutoff scale. Otherwise the nonrenormalizable coupling G would be zero and the dynamical picture would not be valid. In other words, the poles of μ_1 and g_Tg_B respectively μ_2 and g_Eg_N must exactly cancel.

To check this we have to come back to a subtle point: The full one loop renormalization group running is not an appropriate tool to discuss the behaviour around the pole since the loop expansion breaks down in that region. However, as we have already mentioned, it is possible to discuss the pole region in a $1/N_{c'}$ expansion [10]. $N_{c'}$ is not the number of conventional colour degrees of freedom but is connected to the new interaction that produces the bound states. Therefore gluon contributions won't appear in our lowest order $1/N_{c'}$ beta–functions. The expansion parameter $1/N_{c'}$ of course has to apply to the leptonic fields as well as the quark fields since they both play the role of bound state constituents. $N_{c'}$ has an obvious meaning in models where the bound states are produced by a new gauge interaction. In models which try to produce the bound states at the Planck scale by supergravity or stringy effects the meaning and justification of the $1/N_{c'}$ expansion is not that clear and eventually has to arise from the details of such model. In any case the validity of this expansion has to be assumed to have any control over the strong dynamics.

The leading order $1/N_{c'}$ beta functions for our low energy model are:

$$
\begin{aligned}
\frac{d}{dt}\mu &= \frac{1}{16\pi^2}(3g_T^2+3g_B^2)\mu && (15)\\
\frac{d}{dt}\mu' &= \frac{1}{16\pi^2}(g_N^2+g_E^2)\mu' && (16)\\
\frac{d}{dt}g_T &= \frac{1}{16\pi^2}(3g_T^2)g_T && (17)\\
\frac{d}{dt}g_B &= \frac{1}{16\pi^2}(3g_B^2)g_B && (18)\\
\frac{d}{dt}g_N &= \frac{1}{16\pi^2}(g_N^2)g_N && (19)\\
\frac{d}{dt}g_E &= \frac{1}{16\pi^2}(g_E^2)g_E && (20)
\end{aligned}
$$

One easily sees that $d\frac{\mu}{g_T g_B}/dt = d\frac{\mu}{g_N g_E}/dt = 0$. This is the consequence of the fact that in lowest order $1/N_{c'}$ the Landau poles of the Yukawa couplings as well as of the μ–parameters are solely induced by the renormalization of the Higgs propagators.

Next we have to take a look at the soft breaking terms. The low energy effective theory has soft breaking mass terms of the form

$$\int d^2\theta d^2\bar{\theta}\bar{H}_i e^{2V_{H_i}} H_i m_{H_i}^2 \theta^2\bar{\theta}^2 \qquad (21)$$

with $i = 1..4$. We want these terms to vanish after reparameterization at the scale Λ. This translates into avoiding a pole at Λ for the unreparameterized term. Taking H_1 as an example the lowest order $1/N_{c'}$ beta function is

$$\frac{d}{dt}m_{H_1}^2 = \frac{1}{8\pi^2}3g_B^2(m_{H_1}^2+m_B^2+m_T^2+A_1^2) \qquad (22)$$

m_T and m_B do not run in lowest order $1/N_{c'}$. The trilinear soft breaking terms A_i however are driven into a pole by the Yukawa couplings if they don't vanish at Λ. Thus we have the two constituent conditions

$$A(\Lambda) = 0 \tag{23}$$

$$-m^2_{H_1}(\Lambda) = m^2_B(\Lambda) + m^2_T(\Lambda) = 2\Delta^2 \tag{24}$$

where we have used the universal mass Δ^2 introduced in eq.(13). With these conditions fulfilled, F–term induced dynamical electroweak breaking in fact is a viable model.

The overall picture now shows a nonrenormalizable interaction Lagrangian

$$\begin{aligned}\mathcal{L}^{int}_F = G\int d^2\theta QB^cQT^c(1+\delta(\theta^2) \quad &+ \quad G\int d^2\bar{\theta}\bar{Q}\bar{B}\bar{Q}\bar{T}(1+\bar{\delta}(\bar{\theta}^2) \\ +G'\int d^2\theta LE_\tau LN_\tau(1+\delta'(\theta^2) \quad &+ \quad G'\int d^2\bar{\theta}\bar{L}\bar{E}_\tau\bar{L}\bar{N}_\tau(1+\bar{\delta}'(\bar{\theta}^2)\end{aligned} \tag{25}$$

with $\delta = -\hat{B}$, $\delta' = -\hat{B}'$, $G = \frac{\hat{g}_T\hat{g}_B}{\hat{\mu}}$ and $G' = \frac{\hat{g}_N\hat{g}_E}{\hat{\mu}'}$. The assumption of Yukawa unification would imply $G = G'$. The interaction structure of eq.(25) provides dynamical electroweak breaking via a critical self consistence equation. The low energy effective theory is a slight extension of the MSSM with four Higgs doublets and neutrino Yukawa couplings. The parameters of the effective theory are restricted by the constituent conditions eq.(14), eq.(24) and $tan\beta \approx m_t(\Lambda)/m_b(\Lambda)$. The constituent relations for the four Higgs fields are in the simplest case without mixing

$$H_1 = \frac{\hat{g}_T}{\hat{\mu}}QT^c \tag{26}$$

$$H_2 = \frac{\hat{g}_B}{\hat{\mu}}QB^c \tag{27}$$

$$H_3 = \frac{\hat{g}_N}{\hat{\mu}}LN^c \tag{28}$$

$$H_4 = \frac{\hat{g}_E}{\hat{\mu}}LE^c \tag{29}$$

In component fields this corresponds to

$$h_1 = \frac{\hat{g}_T}{\hat{\mu}}\tilde{t}^\dagger\tilde{q}, \ etc. \tag{30}$$

As discussed earlier, in contrary to the D–term case F–term induced dynamical electroweak breaking cannot involve fermion–condensation any more, the scalar boundstates are entirely made of sfermions. One can understand this model as a realization of the idea of having electroweak breaking induced by the scalar superpartners of the SM fermions. This however does not happen via a tree level potential, but via a dynamical mechanism.

Finally one should make some remarks on prospects and experimental testability of the described model. The phenomenology of F–term induced dynamical electroweak breaking below the Planck scale represents a specific choice in the parameter space of a slightly extended MSSM. This set of parameters is similar to that suggested by some Yukawa unification scenarios which favour Yukawa couplings close to the infrared fixed point [12]. Thus a potential discovery of SUSY signatures could contradict or favour but not easily prove the discussed models. It would be interesting to investigate the embedding of F–term induced dynamical electroweak breaking into a GUT scenario, e.g. SO(10) unification. Generally GUT scenarios that unify Yukawa couplings require a non–universal mass pattern for the scalar soft breaking masses at the GUT scale [11, 9]. Since the conditions on the mass parameters are Planck scale conditions, this basic requirement would be fulfilled in our case. To be viable, our model needs a consistent method to construct its enhanced non–renormalizable operators from an underlying theory. The question whether this can be achieved most likely has to be decided in the framework of a strongly coupled string theory.

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