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Status of Supersymmetry Breaking in String Theory*

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ABSTRACT

Recent developments in understanding supersymmetry breaking in superstring theory are summarized.

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Superstring theory appears to be a viable candidate for a consistent theory of quantum gravity which also bears the possibility of unifying of all known interactions.[1] However, in order to extract physical properties of string theory at the weak scale various technical and conceptual problem have to be solved. First, it is presently not understood how supersymmetry can break at a scale much lower than the Planck scale. Such a hierarchical supersymmetry breaking seems to be a necessary ingredient for obtaining the right low energy physics. Secondly, the (unified) string gauge coupling constant is determined by the vacuum expectation value (Vev) of a gauge neutral scalar field called the dilaton S. To all orders in string perturbation theory S turns out to be a flat direction of the effective potential and thus the value of the gauge coupling constant remains undetermined. Thirdly, there is no known mechanism which keeps the cosmological constant at zero after supersymmetry breaking. Finally, there is an enormous degeneracy in string vacua and it is not understood how the theory chooses between these vast number of possibilities. It is believed that all four problems can only be addressed (or solved) when string non-perturbative effects are taken into account. Unfortunately our understanding of the non-perturbative structure of string theory is very limited. Therefore the theory is commonly analyzed under the assumption that the dominant non-perturbative effects are field theoretical in nature. This then leads to a 'superstring-inspired' treatment of non-perturbative effects.

Recently there has been considerable progress in understanding non-perturbative properties of (non-critical) string theories in lower space-time dimensions.[2] In particular it has been argued that 'stringy' non-perturbative effects are larger than their field theory counterparts.[3] If this situation persists for critical string theory any field theoretical treatment of non-perturbative effects in string theory would be invalidated. Here, however, we focus on supersymmetry breaking in a purely field theoretical context.

The prominent non-perturbative effect for breaking supersymmetry dynamically is gaugino condensation.[4] This was already investigated in the context of supergravity in the early 80'ies and was first applied to the heterotic string in ref. 5.

Let us briefly summarize the mechanism proposed in supergravity which explains the hierarchy of the weak scale M_W, the supersymmetry breaking scale M_S and the Planck scale $M_{\rm Pl}$.[4] This is done in two steps. First, imagine that supersymmetry is broken by the Vev of a dimension 2 (auxiliary) field F ($\langle F\rangle \sim M_S$) in a so called 'hidden sector' which has no renormalizable couplings with the observable sector. This breaking is then communicated to the observable sector via gravitational interactions and results in masses m_s for the supersymmetric partners of the observable spectrum of the order $m_s \sim M_S^2/M_{\rm Pl}$. Thus for m_s to be in the TeV range one needs $M_S \sim 10^{10} - 10^{11} GeV$. In a second step (which will occupy the rest of the talk) one explains the generation of $M_{\rm S}$ from $M_{\rm Pl}$. The favorite scenario is to assume the hidden sector to consist of an asymptotically free non-abelian gauge theory which is weakly coupled at $M_{\rm Pl}$ but becomes strongly coupled at some lower scale $\Lambda_{\rm c}$. In such a theory the gauginos λ will condense and possibly induce supersymmetry breaking at a scale $M_{\rm S}^2 \sim \Lambda_{\rm c}^3/M_{\rm Pl}$. In order to achieve the above hierarchy the condensation should occur at $\Lambda_{\rm c} \sim 10^{13} - 10^{14} GeV$. $\Lambda_{\rm c}$ can be estimated by RG-invariance to be[4]

$$\Lambda_{\rm c} \sim M_{\rm Pl} \exp(-\frac{8\pi^2}{b_0 g^2(M_{\rm Pl})}) \qquad (1)$$

where g is the gauge coupling constant and b_0 the one loop coefficient of the β-function. Thus, to achieve the desired hierarchy one needs to adjust the gauge coupling constant and the size of the hidden gauge group.

The scenario can be applied to superstring inspired models. The occurrence of a hidden sector is very generic in string vacua; the hidden E_8 of the heterotic string is only one of many examples. The new ingredient is a dynamical gauge coupling constant which is determined by the Vev of the Dilaton superfield ($1/g^2 = \langle {\rm Re} S\rangle$) at the string tree level. This relation is further corrected at the one loop level by so called threshold corrections Δ which arise from integrating out heavy modes of the string spectrum.[6] If the masses of the heavy fields depend on the Vev of a light field T (Higgs or modulus) the threshold corrections can also be field dependent and we have altogether

$$1/g^2 = \text{Re}S + \Delta(T, \bar{T})\,. \tag{2}$$

Such a field dependent contribution to $1/g^2$ occurs frequently in string theory.[7] Thus, from eqs. (1) and (2) we see that the essential difference in these 'string inspired' scenarios is the fact that the condensation scale is field dependent and undetermined. A self consistent analysis must be invoked which assumes the formation of the condensates, then derives an effective potential for it which indeed leads to a non-zero condensate after minimization.

It has not been entirely clear how to incorporate field dependent gauge couplings into a manifestly supersymmetric gaugino condensation mechanism. Therefore let me first mention some recent investigations about the structure of gauge couplings in supersymmetric theories.[7–10] At the tree level, supersymmetry constrains field dependent gauge couplings to be the real part of a holomorphic function $f(\phi)$. (Its imaginary part is related to a field dependent θ-angle.) This turns out to be no longer true once loop effects are included. In particular the threshold corrections Δ need not be the real part of any holomorphic function. Such non-holomorphic contributions can arise when massless particles contribute in loops to Δ. This situation was found to occur in some superstring models[7] and a general formula for the non-harmonic contribution to Δ in terms of the tree level couplings of the massless fields was given in ref. 10. If there are no massless particles in the theory or if only the massive modes are integrated out $1/g^2$ remains the real part of a holomorphic function also at the loop level. Therefore it is crucial to distinguish between the physical (running) gauge coupling and the Wilsonian gauge coupling g_W.[8] The latter is defined to receive loop contributions only from massive modes and thus continues to be the real part of a holomorphic function. It was further shown that g_W is not renormalized beyond one loop.[8,11]

Now we are prepared to discuss gaugino condensation in a supersymmetric framework including field dependent gauge couplings. Most of the effects of the condensation are captured by a supersymmetric effective Lagrangian below the condensation scale where all gauge degrees of freedom have been integrated out.[12] One introduces a composite chiral superfield U describing the dynamics of the condensate and derives the effective Lagrangian by demanding that the (anomalous) Ward identities of the theory are satisfied. For a hidden sector which contains a pure gauge theory the effective superpotential reads[12,13,8,10]

$$W^{\text{eff}}(U,S,T) = \frac{U}{4}\left[f_W(S,T) + \frac{b_0}{24\pi^2}\ln(\frac{U}{M_{\text{Pl}}^3})\right] \tag{3}$$

where the new feature is the appearance of f_W (which contains the Wilsonian gauge couplings g_W as its real part) in eq. (3). As we just mentioned f_W does not get corrected beyond one loop and hence eq. (3) is exact. Minimization of W^{eff} with respect to U leads to

$$U \sim M_{\text{Pl}}^3 \exp\left(-\frac{24\pi^2}{b_0} f_W(S,T)\right) . \tag{4}$$

We see that U is closely related to the condensation scale Λ_c and as we expected depends on the fields S, T. (The precise relation between Λ_c and U reads $\Lambda_c^3 = e^{K/2}|U|$ where K is the Kähler potential.) Minimization of W with respect to S, however, leads to $S = \infty, U = 0$. Thus the effective potential (3) predicts a vanishing gauge coupling constant, clearly an unacceptable state of affairs. A way out of this dilemma was proposed some time ago in ref. 14 and further investigated in ref. 15. It imagines a hidden sector of two gauge groups with very closely matched β-functions (e.g. $SU(N) \times SU(N+1)$) and a difference δ ($\delta = \frac{\Delta_{N+1}}{N+1} - \frac{\Delta_N}{N}$) in the field independent, thus constant threshold corrections. Under these assumptions minimization of eq. (3) generates a dilaton Vev whose size is controlled by large N and to leading order in N is given by

$$\text{Re}S = \frac{N^2\delta}{16\pi^2} \tag{5}$$

(For example $SU(9) \times SU(10)$ with $\delta = 4$ leads to Re$S = 2.1$ or $\alpha_{\text{GUT}} \sim 1/25$ but a slightly too large $\Lambda_c \sim 10^{16}$GeV.) Unfortunately, without any T dependence of f_W supersymmetry is unbroken at this minimum.

However, once the dilaton is stabilized supersymmetry can be broken in the presence of a T dependent f_W which seems to be generically induced in string models but has been calculated explicitly only for a certain class of string vacua (orbifolds) as a function of the (untwisted) moduli T.[7] In string theory one does not compute f_W directly but rather Δ which includes contributions of the light fields. f_W is obtained after subtracting these contributions of the light fields which leads in the case of the E_8 gauge group in symmetric orbifolds to[10,13]

$$f_W^{E_8} = S + \frac{b_0}{8\pi^2} \ln \eta^2(T) \tag{6}$$

where η is the Dedekind η-function. Inserted into eq. (4) results in

$$U \sim M_{\rm Pl}^3 \, \eta^{-6}(T) \exp(-\frac{24\pi^2 S}{b_0}) \,. \tag{7}$$

It was shown in ref. 13 that the superpotential corresponding to eq. (7) can lead to supersymmetry breaking via $\langle F_T \rangle$ (the auxiliary field in the T supermultiplet) once the dilaton is fixed.

The T-dependence of eq. (7) was first derived from a slightly different point of view. For these orbifold vacua there is a discrete symmetry (modular invariance) acting on the moduli fields T.[16] Ref. 17 investigated the possible T dependence of $W^{\rm eff}$ which is compatible with modular covariance, holomorphicity and regularity of $W^{\rm eff}$. These requirements constrain the T dependence to the one of eq. (7). Since modular invariance is believed to hold to all orders in string perturbation theory, we understand the exactness of this superpotential from both points of view.

Including the threshold corrections into the gaugino condensation mechanism (partially) solves another of the problems listed above. Similar to the dilaton the moduli T are also scalar fields whose Vev is undetermined to all orders in string perturbation theory. They parametrize a family of degenerate string vacua. The non-perturbative superpotential generated via gaugino condensation determines their Vevs thus (partially) lifting the vacuum degeneracy. In the case of orbifolds the untwisted modulus T corresponds to the radius of the internal manifold. It is interesting that gaugino condensation seems to fix the radius at some finite value thus leading to compactification through a dynamical mechanism.[13] Unfortunately, supersymmetry breaking does depend on the functional form of f_W which prohibits a generic analysis for all string vacua.

Finally, let us consider a hidden sector which contains $N_{\rm f}$ generations of (massive) matter fields C^I in the fundamental representation of a hidden gauge group $SU(N)$. This is a more realistic situation in string theory than a pure gauge hidden sector. The effective potential (3) is generalized for this case to read[12,18,10]

$$W^{\rm eff}(U,\Pi,S,T) = \frac{U}{4}\left[f_W(S,T) + \frac{N-N_{\rm f}}{8\pi^2}\ln\frac{U}{M_{\rm Pl}^3} + \frac{1}{8\pi^2}\ln\frac{\det\Pi}{M_{\rm Pl}^{2N_{\rm f}}}\right] + W^{\rm tree}(\Pi,T) \tag{8}$$

where $\Pi^{IJ} = \langle C^I C^J\rangle$ denotes the condensate of the matter fields. The analysis of this superpotential becomes somewhat more involved and depends on the detailed form of $W^{\rm tree}(\Pi,T)$ and in particular on the generation of the mass term. So far only special cases have been studied in detail.[18,10] However one can note a few general features. After integrating out all fields except S and T an equation quite similar to (7) emerges. (Again the T dependence follows from modular invariance.) The mechanism for stabilizing the dilaton can be implemented as before with the possibility of an improved hierarchy and supersymmetry can again be broken by $\langle F_T\rangle$. One should also note that a hidden sector without a mass term but rather a quartic coupling can appear in string theory and it is straightforward to treat such a case in the framework of an effective Lagrangian. Overall, for a hidden sector containing matter representation a more detailed analysis is needed.

Unfortunately, almost all of the above scenarios lead to a non-zero cosmological constant and they offer no explanation why it should vanish.

To summarize, we have sketched a possible mechanism which stabilizes the dilaton at a phenomenologically viable value and leads to the generation of an intermediate scale via gaugino condensation. After including field dependent threshold corrections supersymmetry can be broken. Unfortunately the cosmological constant is non-zero.

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