Gluino condensation in strongly coupled heterotic string theory

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Strongly coupled heterotic $E_8\times E_8$ string theory, compactified to four dimensions on a large Calabi-Yau manifold $\mathbf{X}$, may represent a viable candidate for the description of low-energy particle phenomenology. In this regime, heterotic string theory is adequately described by low-energy M theory on $\mathbf{R}^4\times\mathbf{S}^1/\mathbf{Z}_2\times\mathbf{X}$, with the two E_8's supported at the two boundaries of the world. In this paper we study the effects of gluino condensation as a mechanism for supersymmetry breaking in this M-theory regime. We show that when a gluino condensate forms in M theory, the conditions for unbroken supersymmetry can still be satisfied locally in the orbifold dimension $\mathbf{S}^1/\mathbf{Z}_2$. Supersymmetry is then only broken by the global topology of the orbifold dimension, in a mechanism similar to the Casimir effect. This mechanism leads to a natural hierarchy of scales, and elucidates some aspects of heterotic string theory that might be relevant to the stabilization of moduli and the smallness of the cosmological constant. [S0556-2821(96)04324-X]

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I. INTRODUCTION: PHENOMENOLOGY OF *M* THEORY

Despite its remarkable phenomenological promise [1–3], string theory still leaves unanswered many pressing questions about its contact with the low-energy world. Among the issues that we would certainly want to understand better in a unified theory are the mechanism of supersymmetry breaking with a large hierarchy of scales, the stabilization of moduli, and the smallness of the cosmological constant (for reviews and references, see [4–11]).

Our present understanding of this subject indicates that string theory might be able to identify the right degrees of freedom in which phenomenology can be naturally understood. There are, however, equally strong indications that in the regime directly relevant to phenomenology, the natural degrees of freedom of perturbative string theory are strongly coupled [12–15]. Recently, we have witnessed a revolution that is rapidly changing our understanding of string theory in the strongly coupled regime, leading in many cases to a dual description of the physics in this regime in terms of more natural, weakly coupled degrees of freedom. One can wonder whether these dual descriptions might lead to variables more appropriate for the description of low-energy phenomenology.

In the recent studies of string dualities, at least one such new paradigm may have already appeared [16–18]. Compactification of the strongly coupled heterotic $E_8\times E_8$ string theory on a Calabi-Yau manifold $\mathbf{X}$ is most naturally described by 11-dimensional M theory, compactified to $\mathbf{R}^4$ on a manifold with one extra dimension, $\mathbf{X}\times\mathbf{S}^1/\mathbf{Z}_2$ [16]. This extra dimension, invisible at weak heterotic coupling, is an orbifold dimension, and the total space-time manifold has a boundary with two components. At low energies, the effective description of M theory is in terms of 11-dimensional supergravity, coupled to one E_8 Yang-Mills supermultiplet at each boundary of the world [17].

This picture gives an interesting new twist to the old Kaluza-Klein idea. For a low-energy observer, the world first looks four dimensional. After crossing a certain threshold, the world becomes effectively five dimensional, but the matter sector containing the standard model still lives at a four-dimensional boundary. The bulk of the five-dimensional space-time supports only gravity (as well as other fields coming from the 11-dimensional supergravity multiplet). At the other end of the world, another gauge sector, the other E_8 of the heterotic string theory, is hidden and communicates with the matter of the standard model only gravitationally. Finally, at even higher energies, the observer reaches the compactification scale and sees the additional six dimensions compactified on the Calabi-Yau manifold, and the world becomes 11 dimensional.

This newly understood regime of heterotic string theory seems very attractive phenomenologically. In [18], Witten used this regime to analyze the strongly coupled heterotic string compactified on a large Calabi-Yau manifold, with the four-dimensional grand-unified coupling α_{GUT} acceptably small. A detailed analysis reveals that for such compactifications, unlike in the weakly coupled heterotic string, the strengths of all interactions including gravitational can be naturally unified at the unification scale. In other words, the unacceptable prediction of the size of the Newton constant G_N, as made generically by the weakly coupled heterotic string theory, is alleviated at strong coupling, in the M-theory regime.[1] More recently, other interesting phenomenological implications of this scenario have been studied in some detail in [20] (see also [21]).

Gluino condensate, supersymmetry breaking, and the cosmological constant

The unification of couplings, essentially due to the presence of the extra dimension of the type discussed in [16], can be considered one of the first phenomenological successes of

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[1]For a clear nontechnical exposition of this result, see Sec. 4.3 and Fig. 6 of [19].

M theory. This makes one wonder whether *M* theory has anything to say about the mysteries of hierarchical supersymmetry breaking and the smallness of the cosmological constant.

There are three known mechanisms that can trigger supersymmetry breaking in string theory: world-sheet instantons, space-time instantons, and strong infrared dynamics. Superpotentials generated by instantons have been recently studied in [22], where three-dimensional compactifications of *F* theory were used to gain information about four-dimensional physics.

Gluino condensation in the hidden sector is a representation of the third class. It has been extensively studied as a mechanism of supersymmetry breaking, ever since the pioneering papers in supergravity [23] and in string theory [24,25]; for reviews and references on this subject,[2] see [26,27].

In addition to providing a natural mechanism of supersymmetry breaking, gluino condensation could also be relevant to the cosmological constant problem and the stabilization of moduli in string theory. At a very early stage of the studies of gluino condensates in weakly coupled heterotic string theory, it has been noticed [24,25] that certain terms in the Lagrangian of the ten-dimensional heterotic supergravity conspire in a very particular way, leading to a potential of a very special, ''no-scale'' type, first considered in [32,33]. Potentials of the no-scale type have been argued to break supersymmetry while keeping the cosmological constant naturally zero without fine-tuning. The main problem with this mechanism seems to be the apparent absence of a satisfactory symmetry principle that could explain and protect this particular form of the potential and lead to supersymmetry breaking with zero cosmological constant in the presence of quantum effects.

In this paper, we will study gluino condensation in the strongly coupled, *M*-theory regime of the heterotic string theory, to the lowest nontrivial order in a long-wavelength expansion. We will assume that a gluino condensate develops in the E_8 sector hidden at the other end of the world and will study its consequences for supersymmetry breaking.[3] Our analysis will reveal some unexpected properties of gluino condensation in the heterotic string theory at strong coupling. Since the rest of the paper will be somewhat technical, we summarize our results here:

In the low-energy Lagrangian of *M* theory, we find a ''conspiration of terms'' similar to the one observed in the low-energy Lagrangian of the weakly coupled heterotic string theory. Incidentally, this explains some rather singular terms encountered in [17] in the construction of the low-energy effective Lagrangian of *M* theory on a manifold with boundary. It also suggests that when a gluino condensate develops at the boundary, the field strength *G* of the three-form *C* from the 11-dimensional supergravity multiplet develops a compensating vacuum expectation value supported at the boundary. We also encounter first indications that the 11-dimensional variables of *M* theory might be more appropriate for the description of supersymmetry breaking by the gluino condensate—the role of the would-be Goldstino is played by the normal component ψ_{11} of the 11-dimensional gravitino.

Even in the presence of the gluino condensate, we will still be able to solve the unbroken supersymmetry conditions in any given coordinate system. This phenomenon might come as a surprise and is intimately related to the existence of the extra dimension in *M* theory.

The solution of unbroken supersymmetry conditions exists locally, but not globally in the extra dimension. When we try to extend the local solution globally over $\mathbf{S}^1/\mathbf{Z}_2\times\mathbf{X}$, we encounter a topological obstruction (essentially, the total cohomology class of the gluino condensate). Therefore, supersymmetry is broken by the global topology of the extra, orbifold dimension, in a process similar to the Casimir effect.

The fact that the unbroken supersymmetry conditions can be satisfied in any coordinate system on $\mathbf{R}^4\times\mathbf{S}^1/\mathbf{Z}_2\times\mathbf{X}$ leads to an intriguing refinement of the phenomenology of supersymmetry breaking in these models. We have argued that in the *M*-theory scenario, observers at intermediate energies will see the world as five dimensional. At length scales larger than the Calabi-Yau compactification radius, but still much smaller than the radius of the fifth dimension, these observers should see unbroken supersymmetry, even if they are directly at the other end of the world where the gluino condensate has formed. As the resolution is diminished, the supersymmetry breaking effects, caused by the compactness of the fifth dimension, become visible, and at even larger length scales, the world will be effectively four dimensional and supersymmetry will be broken. However, this breakdown should be rather mild, since it is only caused by effects sensitive to the global topology of the fifth dimension.

This mechanism of supersymmetry breaking generates a natural hierarchy of scales. In the phenomenologically interesting regime, the radius of the fifth dimension can be expected [20] to be at least an order of magnitude larger than the 11-dimensional Planck length. The mass of the five-dimensional gravitino is only induced quantum mechanically by loop effects sensitive to the size of the fifth dimension and is therefore suppressed by a power of the inverse radius of the fifth dimension.

This hidden 11-dimensional supersymmetry, broken only by the global topology of the orbifold dimension, explains the ''conspiracy'' that leads in the weakly coupled heterotic string theory to the no-scale potential with supersymmetry breaking and zero cosmological constant at the tree level.

[2]The simplest version of gluino condensation in weakly coupled heterotic string theory does not successfully explain the hierarchy of scales and the stabilization of moduli. Several suggestions how to alleviate these problems without abandoning the regime of weak string coupling have been made [3], among them the racetrack models with multiple gluino condensates [28,29] or modified gauge-kinetic functions [30,31].

[3]The formation of a gluino condensate is exactly what one expects on the basis of a simple physical argument. For the compactifications studied in [18] and in the present paper, strong coupling in the hidden E_8 develops exactly when the other couplings attain phenomenologically interesting values. With strong gauge coupling in the hidden sector, we can expect that a gluino condensate is dynamically generated. This aspect of the strong gauge dynamics should be stable under the effects caused by the coupling to gravity.

II. GLUINO CONDENSATION IN HETEROTIC STRING THEORY AND *M* THEORY

In this paper we will study the heterotic $E_8\times E_8$ string theory compactified to $\mathbf{R}^4$ on a Calabi-Yau threefold $\mathbf{X}$. In the strong coupling limit, the low-energy description of this theory is in terms of 11-dimensional M theory compactified on $\mathbf{R}^4\times\mathbf{S}^1/\mathbf{Z}_2\times\mathbf{X}$, with the two E_8 gauge groups supported at the two space-time boundaries in the orbifold dimension $\mathbf{S}^1/\mathbf{Z}_2$. We will deal with various supergravities that describe the low-energy physics of such compactifications.

Our 10-dimensional and 11-dimensional conventions are as in [16]. The space-time signature is $-+\cdots+$. Eleven-dimensional vector indices will be written as $I,J,K,\ldots$. The 11-dimensional Γ matrices are 32×32 real matrices satisfying $\{\Gamma_I,\Gamma_J\}=2g_{IJ}$, with $g_{IJ}=\eta_{mn}e_I^{\ m}e_J^{\ n}$ the 11-dimensional metric. Each of the two boundary components of the 11-dimensional manifold supports one E_8 Yang-Mills supermultiplet. One of the E_8's will be broken by the spin connection embedding to a grand-unified E_6 group, while the other E_8 will be strongly coupled and hidden at the other end of the world. The adjoint index of this hidden E_8 will be denoted by $a,b,\ldots$.

On $\mathbf{R}^4\times\mathbf{S}^1/\mathbf{Z}_2\times\mathbf{X}$, we will use four-dimensional vector indices $\mu,\nu,\ldots$, which parametrize the flat Minkowski space $\mathbf{R}^4$, and vector indices $i,j,k,\ldots$ and their complex conjugates $\bar{i},\bar{j},\bar{k},\ldots$, which correspond to a complex coordinate system on the Calabi-Yau threefold $\mathbf{X}$. The ten-dimensional vector indices that parametrize $\mathbf{R}^4\times\mathbf{X}$ will be written as $A,B,C,\ldots$. Our other conventions on $\mathbf{X}\times\mathbf{S}^1/\mathbf{Z}_2$ are as in [18].

A. Gluino condensation and the potential at weak coupling

First, we recall some aspects of the gluino condensation in the hidden sector of the weakly coupled heterotic string theory that will be relevant for our purposes.

Consider, as in [24], the weakly coupled heterotic $E_8\times E_8$ string theory compactified on a Calabi-Yau threefold $\mathbf{X}$. On any given $\mathbf{X}$, we have a covariantly constant holomorphic three-form ϵ_{ijk} (and its antiholomorphic complex conjugate $\overline{\epsilon_{ijk}}=\epsilon_{\bar{i}\bar{j}\bar{k}}$). In ten dimensions, $\overline{\chi^a}\Gamma_{ABC}\chi^a$ is the only gluino bilinear that is not identically zero by Fermi statistics and chirality. If this bilinear develops a nonzero vacuum expectation value proportional to the covariantly constant holomorphic three-form on $\mathbf{X}$,

$$\langle\overline{\chi^a}\Gamma_{ijk}\chi^a\rangle=c\Lambda_{E_8}^3\epsilon_{ijk} \tag{2.1}$$

(and similarly for the complex conjugate), the four-dimensional observer will interpret this expectation value as a nonzero gluino condensate $\langle\overline{\chi^a}\chi^a\rangle$ [and $\langle\overline{\chi^a}\gamma_5\chi^a\rangle$]. In Eq. (2.1), Λ_{E_8} is the characteristic scale of the hidden gauge sector, at which the gauge coupling becomes strong, and c is a (complex) number of order 1.

In the process of analyzing the physics of the gluino condensate in weakly coupled heterotic string theory, it has been noticed [24] (see also [2], Vol. 2, p. 326) that the Lagrangian of ten-dimensional heterotic supergravity exhibits a special feature that could lead, at least at the tree level, to supersymmetry breaking with zero cosmological constant without fine-tuning. The argument is roughly as follows. The Lagrangian contains a gluino self-interaction term which is quartic in χ^a; it also contains an interaction between the gluino bilinear $\overline{\chi^a}\Gamma_{ABC}\chi^a$ and the three-form field strength H_{ABC}. Together with the kinetic term H^2, these terms conspire in such a way that they can be assembled into a perfect square,

$$-\frac{3\kappa_{10}^2}{4\lambda_{10}^4}\int_{M^{10}}d^{10}x\sqrt{g}\,\frac{1}{\phi^{3/2}}\,[H_{ABC}-\lambda_{10}^2\sqrt{2}\,\phi^{3/4}(\overline{\chi^a}\Gamma_{ABC}\chi^a)]^2. \tag{2.2}$$

(We have used the normalizations of [24]; ϕ is the ten-dimensional dilaton, while κ_{10} and λ_{10} denote the ten-dimensional gravitational and gauge coupling, respectively.) Consider now the situation in which a gluino condensate has formed, proportional to the covariantly constant three-form on $\mathbf{X}$ as in Eq. (2.1). If we assume that the three-form field strength H_{ABC} develops a compensating vacuum expectation value,

$$H_{ijk}=c\Lambda_{E_8}^3\lambda_{10}^2\sqrt{2}\,\phi^{3/4}\epsilon_{ijk}, \tag{2.3}$$

such that the perfect square term (2.2) in the potential vanishes, the cosmological constant will be zero at the tree level. At the same time, one can show that supersymmetry is broken by the condensates (2.1) and (2.3).

The easiest way to see supersymmetry breaking in the presence of the condensates is to look at the relevant part of the supersymmetry variation of the fermions. There are two relevant fermions in the theory: the ten-dimensional gravitino Ψ_A and the dilatino λ. Schematically, the relevant parts of their supersymmetry variations are given by

$$\begin{aligned}\delta\Psi_A=&\frac{1}{\kappa_{10}}D_A\eta+\frac{\sqrt{2}}{32}\left(\frac{\kappa_{10}}{\lambda_{10}^2}\right)\frac{1}{\phi^{3/4}}H_{BCD}\\&\times(\Gamma_A^{\ BCD}-9\,\delta_A^B\Gamma^{CD})\,\eta-\frac{1}{256}\,\kappa_{10}(\overline{\chi^a}\Gamma_{BCD}\chi^a)\\&\times(\Gamma_A^{\ BCD}-5\,\delta_A^B\Gamma^{CD})\,\eta+\cdots,\\ \delta\lambda=&\cdots+\frac{1}{8}\left(\frac{\kappa_{10}}{\lambda_{10}^2}\right)\frac{1}{\phi^{3/4}}H_{ABC}\Gamma^{ABC}\eta\\&+\frac{\sqrt{2}}{384}\,\kappa_{10}(\overline{\chi^a}\Gamma_{ABC}\chi^a)\Gamma^{ABC}\eta+\cdots.\end{aligned} \tag{2.4}$$

(Here the ellipsis corresponds to terms that are either proportional to the gravitino and dilatino or contain the space-time derivative of the dilaton.) We can see from Eq. (2.4) that in the presence of the condensates (2.1) and (2.3), the unbroken supersymmetry conditions $\delta\Psi_A=0$ and $\delta\lambda=0$ cannot be satisfied. A particular linear combination of Ψ_A and λ behaves as a Goldstone fermion and gives a nonzero tree-level mass to the gravitino, and supersymmetry is broken in this approximation.

In four dimensions, the perfect square structure (2.2) of the heterotic supergravity Lagrangian leads to the superpotential and Kähler potential of the very special, no-scale type [32,33]. Superpotentials and Kähler potentials of the no-scale type were proposed [32] in earlier attempts to link supersymmetry breaking with the solution of the cosmological constant problem (see also the discussion in [11]). One of the main drawbacks of this approach so far has been the apparent lack of a symmetry principle that could explain and protect this particular form of the potential.[4]

B. Strong coupling and *M* theory

At strong string coupling and large radius of the Calabi-Yau manifold, the compactification is effectively described by low-energy M theory on $\mathbf{R}^4\times\mathbf{S}^1/\mathbf{Z}_2\times\mathbf{X}$. The effective Lagrangian for this theory has been constructed in [17]. It contains the 11-dimensional supergravity multiplet $e_I{}^m$, ψ_J, and C_{IJK} in the bulk, coupled to one E_8 Yang-Mills supermultiplet A_B^a, χ^a at each of the two ten-dimensional boundaries.

To order $\kappa^{2/3}$, the Lagrangian is given by

$$\mathcal{L}=\frac{1}{\kappa^2}\int_{M^{11}}d^{11}x\sqrt{g}\Bigg[-\frac{1}{2}R-\frac{1}{2}\overline{\psi_I}\Gamma^{IJK}D_J\left(\frac{\Omega+\hat{\Omega}}{2}\right)\psi_K-\frac{1}{48}G_{IJKL}G^{IJKL}-\frac{\sqrt{2}}{384}(\overline{\psi_I}\Gamma^{IJKLMN}\psi_N+12\overline{\psi^J}\Gamma^{KL}\psi^M)(G_{JKLM}$$
$$+\hat{G}_{JKLM})-\frac{\sqrt{2}}{3456}\epsilon^{I_1I_2\cdots I_{11}}C_{I_1I_2I_3}G_{I_4\cdots I_7}G_{I_8\cdots I_{11}}\Bigg]+\frac{1}{2\pi(4\pi\kappa^2)^{2/3}}\int_{M^{10}}d^{10}x\sqrt{g}\Bigg(-\frac{1}{4}F^a_{AB}F^{aAB}-\frac{1}{2}\overline{\chi^a}\Gamma^A D_A(\hat{\Omega})\chi^a$$
$$-\frac{1}{8}\overline{\psi_A}\Gamma^{BC}\Gamma^A(F^a_{BC}+\hat{F}^a_{BC})\chi^a+\frac{\sqrt{2}}{48}(\overline{\chi^a}\Gamma^{ABC}\chi^a)\hat{G}_{ABC11}\Bigg). \tag{2.5}$$

(Explicit expressions for the supercovariant objects $\hat{\Omega}$, $\hat{F}^a_{AB}$, and $\hat{G}_{IJKL}$ can be found in [17].) The fields of the bulk supergravity multiplet satisfy natural orbifold boundary conditions, discussed in detail in [17]. It was also shown in [17] that the four-form field strength G_{IJKL} satisfies a modified Bianchi identity

$$dG_{11ABC}=-\frac{3\sqrt{2}}{2\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})(\mathrm{tr}F_{[AB}F_{CD]}-\tfrac{1}{2}R_{[AB}R_{CD]}), \tag{2.6}$$

which will be important later in the paper.

The effective Lagrangian (2.5) is invariant under local supersymmetry, whose parameter η satisfies the orbifold condition $\eta(-x^{11})=\Gamma_{11}\eta(x^{11})$. For the purposes of this paper, we will only need the rules for the supersymmetry transformations of the fermions: the relevant supersymmetry transformations are

$$\delta\psi_A=D_A\eta+\frac{\sqrt{2}}{288}G_{IJKL}(\Gamma_A{}^{IJKL}-8\delta^I_A\Gamma^{JKL})\eta$$
$$-\frac{1}{576\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})(\overline{\chi^a}\Gamma_{BCD}\chi^a)$$
$$\times(\Gamma_A{}^{BCD}-6\delta^B_A\Gamma^{CD})\eta+\cdots, \tag{2.7}$$

$$\delta\psi_{11}=D_{11}\eta+\frac{\sqrt{2}}{288}G_{IJKL}(\Gamma_{11}{}^{IJKL}-8\delta^I_{11}\Gamma^{JKL})\eta$$
$$+\frac{1}{576\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})(\overline{\chi^a}\Gamma_{ABC}\chi^a)\Gamma^{ABC}\eta+\cdots,$$
$$\delta\chi^a=-\tfrac{1}{4}F^a_{AB}\Gamma^{AB}\eta+\cdots.$$

The ellipsis denotes terms of order $\kappa^{4/3}$, as well as known terms of order $\kappa^{2/3}$ bilinear in the gravitinos that we will not need.

As we recalled in the previous subsection, the effective supergravity Lagrangian of the weakly coupled ten-dimensional heterotic string theory describes the interaction between the gluino bilinears $\overline{\chi^a}\Gamma_{ABC}\chi^a$ and the three-form field strength H_{ABC}, by the perfect square term (2.2), leading to the no-scale potential and the corresponding mechanism of supersymmetry breaking. At first, one would not expect such a perfect square structure to also appear in the effective Lagrangian of M theory. Indeed, the gluinos of M theory live at the space-time boundary and can only contribute to the Lagrangian through surface terms. On the other hand, the three-form C_{IJK}, whose field strength four-form G_{IJKL} is the M-theory counterpart of the heterotic field strength H_{ABC}, belongs to the supergravity multiplet, and its kinetic term is supported by the bulk of the 11-dimensional manifold. We have indeed seen that the effective Lagrangian (2.5) contains the corresponding terms,

$$-\frac{1}{12\kappa^2}\int_{M^{11}}d^{11}x\sqrt{g}G^2_{ABC11}$$
$$+\frac{\sqrt{2}}{24(4\pi)^{5/3}\kappa^{4/3}}\int_{M^{10}}d^{10}x\sqrt{g}G_{ABC11}(\overline{\chi^a}\Gamma^{ABC}\chi^a). \tag{2.8}$$

Nevertheless, it is intriguing to notice that, in fact, the perfect square structure of the interaction between the gluinos and the bosonic field strength persists also in M theory.

[4]Recently, some attempts have been made [30] to substantiate the no-scale potentials using S duality.

In the construction of the Lagrangian [17], an unusual boundary interaction term was encountered. This term appears at relative order $\kappa^{4/3}$, is quartic in the gluinos, and, most importantly, is proportional to the boundary delta function $\delta(x^{11})$ evaluated at zero:

$$-\frac{\delta(0)}{96(4\pi)^{10/3}\kappa^{2/3}}\int_{M^{10}} d^{10}x\sqrt{g}(\overline{\chi^a}\Gamma_{ABC}\chi^a)^2. \quad (2.9)$$

In [17], the presence of this term in the effective Lagrangian has been inferred from the requirement of local supersymmetry. That argument was rather formal and involved cancellations of infinities. Still, it is interesting that this term turned out [17] with precisely the right coefficient so that it can be combined with the two terms in Eq. (2.8) into a perfect square:

$$-\frac{1}{12\kappa^2}\int_{M^{11}} d^{11}x\sqrt{g}G_{ABC11}^2+\frac{\sqrt{2}}{24(4\pi)^{5/3}\kappa^{4/3}}\int_{M^{10}} d^{10}x\sqrt{g}G_{ABC11}(\overline{\chi^a}\Gamma^{ABC}\chi^a)-\frac{\delta(0)}{96(4\pi)^{10/3}\kappa^{2/3}}\int_{M^{10}} d^{10}x\sqrt{g}(\overline{\chi^a}\Gamma_{ABC}\chi^a)^2$$
$$=-\frac{1}{12\kappa^2}\int_{M^{11}} d^{11}x\sqrt{g}\left(G_{ABC11}-\frac{\sqrt{2}}{16\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})\overline{\chi^a}\Gamma_{ABC}\chi^a\right)^2. \quad (2.10)$$

Of course, we can turn this argument around and claim that the perfect square structure of the Lagrangian provides a rationale for the existence of the rather singular term (2.9) in the effective Lagrangian of [17]. This statement can be given the following more precise meaning. Inspired by the perfect square structure of the Lagrangian as found in Eq. (2.10), we can reassemble terms in the Lagrangian and redefine the fields, so that the Lagrangian and the supersymmetry transformations no longer contain any explicit terms proportional to infinite coefficients such as $\delta(0)$. In what follows, we will shift the field strength four-form G_{IJKL} by a term supported at the boundary and bilinear in the gluinos, and define a modified field strength $\widetilde{G}_{IJKL}$ by

$$\widetilde{G}_{ABC11}=G_{ABC11}-\frac{\sqrt{2}}{16\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})\overline{\chi^a}\Gamma_{ABC}\chi^a,$$
$$\widetilde{G}_{ABCD}=G_{ABCD}. \quad (2.11)$$

This set of redefined fields is probably better suited for the description of the physics at the relevant scales, since it makes the effective Lagrangian free of formal infinities, to the order to which the low-energy field theory was claimed to make sense in [17].

In the next section we will be interested in configurations on $\mathbf{R}^4\times\mathbf{S}^1/\mathbf{Z}_2\times\mathbf{X}$ that preserve four-dimensional Poincaré invariance. In those cases, all components $\widetilde{G}_{\mu JKL}$, with μ the vector index on $\mathbf{R}^4$, will vanish. The equations of motion for the nonzero components of the modified field strength on $\mathbf{R}^4\times\mathbf{X}\times\mathbf{S}^1/\mathbf{Z}_2$ are then

$$D_I\widetilde{G}^{IJKL}=0; \quad (2.12)$$

i.e., they formally coincide with the equations of motion for the unmodified field strength G_{IJKL} in the absence of the gluino condensate. Of course, this fact depends crucially on the perfect square structure of the Lagrangian.

The field strength G_{IJKL} of the three-form C has to satisfy the Bianchi identity (2.6). In the transformed variables, the Bianchi identity becomes

$$d\widetilde{G}_{11ABCD}=-\frac{3\sqrt{2}}{2\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})(\mathrm{tr}F_{[AB}F_{CD]}-\tfrac{1}{2}R_{[AB}R_{CD]})+\frac{\sqrt{2}}{4\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})\partial_{[A}(\overline{\chi^a}\Gamma_{BCD]}\chi^a). \quad (2.13)$$

For a covariantly constant gluino condensate, such as the one in Eq. (2.1), proportional to the covariantly constant holomorphic three-form ϵ_{ijk} on $\mathbf{X}$, the last term in Eq. (2.13) vanishes identically. The Bianchi identity then formally coincides with the Bianchi identity for the unmodified field strength G_{IJKL} in the absence of the gluino condensate.

III. GLUINO CONDENSATE AND SUPERSYMMETRY IN *M* THEORY

Now we would like to solve the equations of motion in the presence of the gluino condensate. Our strategy will be as follows. First, we find a solution of the equations of motion and the Bianchi identity for the four-form $\widetilde{G}_{IJKL}$. Then we will try to solve the conditions for unbroken supersymmetry, *a priori* expecting an obstruction that should prevent us from finding unbroken supersymmetry in the presence of a gluino condensate. It will come as a surprise that, because of the presence of the extra orbifold dimension of M theory, the unbroken supersymmetry conditions *can* actually be satisfied, locally in the extra dimension. The expected obstruction will only be topological in nature and will prevent us from extending the local solution globally over the extra dimension.

As a first step, we have to solve the equations of motion and the Bianchi identity for the four-form field strength, which in the presence of a covariantly constant gluino condensate on $\mathbf{X}$ are

$$D^I \widetilde{G}_{IJKL} = 0,$$

$$d\widetilde{G}_{ABCD11} = -\frac{3\sqrt{2}}{2\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})(\mathrm{tr}F_{[AB}F_{CD]} - \tfrac{1}{2}R_{[AB}R_{CD]}) - \frac{3\sqrt{2}}{2\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11}-R_{11}) \times (\mathrm{tr}F_{[AB}F_{CD]} - \tfrac{1}{2}R_{[AB}R_{CD]}). \quad (3.1)$$

(Here we have explicitly included, unlike in our previous discussion, the contribution from the other boundary, located at $x^{11}=R_{11}$.)

The Bianchi identity in Eqs. (3.1) cannot be satisfied unless the total cohomology class of its right-hand side vanishes,

$$\sum [F\wedge F] - [R\wedge R] = 0. \quad (3.2)$$

In the compactifications most directly relevant to phenomenology, this condition is satisfied by embedding the spin connection into one of the gauge groups, which is then broken from E_8 to E_6. This embedding makes $\mathrm{tr}\, F\wedge F - R\wedge R$ vanish pointwise in the Calabi-Yau manifold, but does not make the right-hand side of the Bianchi identity in Eqs. (3.1) zero pointwise. As argued in [18], this generates a gradient for the four-form field strength, which is therefore generically nonzero in this particular class of M-theory compactifications.[5] Since the source of $d\widetilde{G}$ is of order $\kappa^{2/3}$ in the long-wavelength expansion, $\widetilde{G}_{IJKL}$ will also be of order $\kappa^{2/3}$.

We have seen in the previous section that Eq. (3.1) for $\widetilde{G}_{IJKL}$ in the presence of a covariantly constant condensate coincide with the equations for the unmodified four-form G_{IJKL} in the absence of the condensate. These equations have been solved, to the same order in $\kappa^{2/3}$ that we are interested in, by Witten in [18]. To solve Eqs. (3.1), we can take any solution $\mathcal{G}_{IJKL}$ from [18] and set

$$\widetilde{G}_{IJKL} = \mathcal{G}_{IJKL}. \quad (3.3)$$

Notice that in accordance with the argument presented at the end of the previous subsection, it is indeed the modified field strength $\widetilde{G}_{IJKL}$, rather than the original G_{IJKL}, that is better behaved near the boundary in the presence of the gluino condensate. In particular, when a gluino condensate $\langle\overline{\chi^a}\Gamma_{ABC}\chi^a\rangle$ forms at the boundary, $\widetilde{G}_{IJKL}$ stays finite and continuous in the vicinity of the boundary, while the original field strength G_{IJKL} develops a rather singular, compensating vacuum expectation value supported at the boundary, $G_{ABC11}\sim\delta(x^{11})\langle\overline{\chi^a}\Gamma_{ABC}\chi^a\rangle$.

The next step is to look at the unbroken supersymmetry conditions $\delta\psi_A=\delta\psi_{11}=\delta\chi^a=0$, with the supersymmetry variations $\delta\psi_A$, $\delta\psi_{11}$, and $\delta\chi^a$ given by Eqs. (2.7). The gluino condensate is of order 1 at the scale where the strong coupling develops in the Yang-Mills sector; therefore, the contribution of the gluino condensate to the supersymmetry variations (2.7) is of order $\kappa^{2/3}$. On the other hand, $\widetilde{G}_{IJKL}$ contributes already at order κ^0, but since it only acquires nonzero values of order $\kappa^{2/3}$, both effects are of the same order in the long-wavelength expansion in powers of $\kappa^{2/3}$.

In terms of the redefined fields, the supersymmetry variations (2.7) take the interesting form[6]

$$\delta\psi_A = D_A\eta + \frac{\sqrt{2}}{288}\widetilde{G}_{IJKL}(\Gamma_A{}^{IJKL} - 8\,\delta_A^I\Gamma^{JKL})\,\eta + \cdots,$$

$$\delta\psi_{11} = D_{11}\eta + \frac{\sqrt{2}}{288}\widetilde{G}_{IJKL}(\Gamma_{11}{}^{IJKL} - 8\,\delta_{11}^I\Gamma^{JKL})\,\eta + \frac{1}{192\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})(\overline{\chi^a}\Gamma_{ABC}\chi^a)\Gamma^{ABC}\eta + \cdots. \quad (3.4)$$

Here the ellipsis again denotes terms of order $\kappa^{4/3}$.

Two aspects of these formulas are worth pointing out.

(1) The gluino condensate drops out from the supersymmetry variation of ψ_A, and it is therefore the normal component ψ_{11} of the 11-dimensional gravitino that plays the role of the would-be Goldstino in the theory. This indicates that the variables of M theory are perhaps better suited for the description of the super-Higgs effect in the heterotic string than those of the weakly coupled theory.

(2) In the supersymmetry variation of ψ_{11}, the term bilinear in the gluinos is accompanied by a term that depends on the normal derivative of the spinor, $D_{11}\eta$.

These two facts represent yet another ''conspiracy'' in the microscopic Lagrangian of M theory on the manifold with boundary and will be crucial in our subsequent analysis of supersymmetry breaking in the presence of the gluino condensate. In particular, this ''conspiracy'' will allow us to solve the unbroken supersymmetry conditions in the vicinity of the boundary where the gluino condensate forms. Indeed, with the gluino condensate appearing only in the condition for the vanishing of $\delta\psi_{11}$, where the $D_{11}\eta$ term appears, one can now hope to solve these conditions by allowing η to depend on x^{11} appropriately. This is to be contrasted with the analogous situation in the theory dimensionally reduced to ten dimensions, which corresponds to the weakly coupled heterotic theory. In the dimensionally reduced theory, the $D_{11}\eta$ term will be absent from $\delta\psi_{11}$, and supersymmetry will necessarily be broken by the gluino condensate in this approximation.

A. Local solution of the unbroken supersymmetry conditions

In the previous section we have noticed that both the equations of motion and the Bianchi identity of the modified field strength $\widetilde{G}_{IJKL}$ coincide with the equations for the unmodified G_{IJKL} in the absence of the condensate and can therefore be solved using the results of [18]. Since $\delta\psi_A$ was also shown to be independent of the gluino condensate, we

[5] This fact could be relevant to the stabilization of moduli in M theory.

[6] Since there are no corrections at this order in $\kappa^{2/3}$ to $\delta\chi^a$ [17], the corresponding equation $\delta\chi^a=0$ is solved at order κ^0 just as in the Calabi-Yau compactifications at weak coupling, and we will drop it from now on.

can extend this argument and start with any solution of the system of equations studied in [18] and use it directly to solve our equations *in the presence* of the condensate.

A solution of the unbroken supersymmetry conditions in the absence of the gluino condensate to order $\kappa^{2/3}$ is represented [18] by a four-form field strength $\mathcal{G}_{IJKL}$ of order $\kappa^{2/3}$ (which we set equal to our modified field strength $\widetilde{G}_{IJKL}$), a metric on $\mathbf{X}\times\mathbf{S}^1/\mathbf{Z}_2$ (which differs by effects of order $\kappa^{2/3}$ from the product of the Ricci-flat metric on $\mathbf{X}$ and the canonical metric on $\mathbf{S}^1/\mathbf{Z}_2$), and a spinor η (which differs from the covariantly constant spinor η_0 on $\mathbf{X}$ by terms of order $\kappa^{2/3}$). The existence of such a solution in the absence of the gluino condensate has been shown in [18].

The formation of a gluino condensate is also an effect of order $\kappa^{2/3}$ and will further modify $\widetilde{\eta}$. On the other hand, since the gluino condensate decouples in our modified variables from all equations except $\delta\psi_{11}=0$, the four-form field strength and the metric will not be modified by the presence of the condensate.

To find a solution of the unbroken supersymmetry conditions in the presence of the gluino condensate to order $\kappa^{2/3}$, the last equation that remains to be satisfied is $\delta\psi_{11}=0$ or, more explicitly,

$$D_{11}\eta' = -\frac{\sqrt{2}}{288}\widetilde{G}_{IJKL}(\Gamma_{11}{}^{IJKL} - 8\,\delta^{I}_{11}\Gamma^{JKL})\,\eta' - \frac{1}{192\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})(\overline{\chi^a}\Gamma_{ABC}\chi^a)\Gamma^{ABC}\eta' + \cdots , \tag{3.5}$$

with the ellipsis again denoting higher order terms in $\kappa^{2/3}$.

Given that $\widetilde{\eta}$ solves Eq. (3.5) in the absence of the gluino condensate, the equation to be actually solved at order $\kappa^{2/3}$ is

$$\partial_{11}(\eta' - \widetilde{\eta}) = -\frac{1}{192\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})(\overline{\chi^a}\Gamma_{ABC}\chi^a)\Gamma^{ABC}\eta_0 . \tag{3.6}$$

This equation has a very simple solution:

$$\eta' = \widetilde{\eta} - \frac{1}{384\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\epsilon(x^{11})(\overline{\chi^a}\Gamma_{ABC}\chi^a)\Gamma^{ABC}\eta_0 . \tag{3.7}$$

This spinor η', which differs from $\widetilde{\eta}$ and therefore from the covariantly constant spinor η_0 on $\mathbf{X}$ by terms of order $\kappa^{2/3}$, thus satisfies the last of the unbroken supersymmetry conditions, Eq. (3.5), in the vicinity of the gluino condensate to the required order in $\kappa^{2/3}$.

Of course, we have to check that η' still satisfies the rest of the unbroken supersymmetry conditions $\delta\psi_A=0$. This is indeed the case to order $\kappa^{2/3}$, as the gluino condensate is covariantly constant. Also, for this spinor to be well defined on the 11-dimensional orbifold, it has to be even under the $\mathbf{Z}_2$ action that defines the orbifold:

$$\eta'(-x^{11}) = \Gamma_{11}\eta'(x^{11}). \tag{3.8}$$

The η' of Eq. (3.7) indeed satisfies this chirality condition, in an interesting way. While $\widetilde{\eta}$ is chiral in ten dimensions and satisfies $\Gamma_{11}\widetilde{\eta}=\widetilde{\eta}$, the second term in Eq. (3.7) is proportional to $\Gamma_{ABC}\widetilde{\eta}$ which is antichiral in ten dimensions,[7] $\Gamma_{11}\Gamma_{ABC}\widetilde{\eta}=-\Gamma_{ABC}\widetilde{\eta}$. In η', this antichiral spinor is, however, multiplied by the step function $\epsilon(x^{11})$ which is odd under the change of orientation of the 11th dimension $x^{11}\to -x^{11}$. Thus η' is even under the combined action of ten-dimensional chirality and orientation reversal of the 11th dimension, and satisfies the orbifold condition (3.8). Hence, surprisingly enough, the presence of the 11th dimension of *M* theory has allowed us to solve the unbroken supersymmetry conditions in the vicinity of the space-time boundary that supports the gluino condensate.

So far, we have not taken into account the global topology of the orbifold dimension. Strictly speaking, our analysis therefore shows that in the presence of the gluino condensate, supersymmetry is unbroken in the formal limit of infinitely strong heterotic string coupling, i.e., as we send R_{11} to infinity. In this limit, η' of Eq. (3.7) would be a globally well-defined solution of the unbroken supersymmetry conditions, to order $\kappa^{2/3}$.

B. Global obstructions and supersymmetry breaking

So far, we have seen that even the observer located directly at the boundary where the gluino condensate forms will see unbroken supersymmetry, as long as the other, weakly coupled boundary is far away. Now we will try to extend the local solution (3.7) of the unbroken supersymmetry conditions to a global solution defined everywhere in $\mathbf{R}^4\times\mathbf{S}^1/\mathbf{Z}_2\times\mathbf{X}$, for a finite radius of the orbifold dimension.

When we try to do so, we encounter an obstruction. We have already solved the unbroken supersymmetry conditions at the end with the strongly coupled E_8 sector, where the gluino condensate forms. The unbroken supersymmetry conditions are also satisfied everywhere in the bulk, and so they only remain to be satisfied at the weakly coupled E_6 end. Since there is no gluino condensate at this weakly coupled end, the unbroken supersymmetry conditions simply require η' to be continuous across this boundary:

$$\eta'(-R_{11}) = \eta'(R_{11}). \tag{3.9}$$

However, the chirality properties of η' discussed in the previous subsection can be used to show that the condition (3.9) is violated if the gluino condensate at the strongly coupled end is nonzero. Indeed, while $\widetilde{\eta}$ is even under $x^{11}\to -x^{11}$, the term proportional to $\epsilon(x^{11})$ in η' is odd under this transformation. Therefore, a topological obstruction must exist that breaks supersymmetry globally, even though we can solve the unbroken supersymmetry conditions locally in any chosen coordinate system.

Now we would like to understand more precisely the nature of this topological obstruction. To do so, it is natural to consider a slightly more general case, in which gluino condensates are allowed to form at both boundaries of the space-time manifold.

Notice first that the gluino condensate $\langle\overline{\chi^a}\Gamma_{ABC}\chi^a\rangle$ is proportional to the components of a three-form on $\mathbf{X}$, but it is

[7] Thus the spinor that represents the unbroken supersymmetry does not have a definite ten-dimensional chirality; however, it still satisfies the chirality condition in four dimensions, $\gamma_5\eta'=\eta'$.

actually better to think of it as a four-form on $\mathbf{X}\times\mathbf{S}^1/\mathbf{Z}_2$. Indeed, the δ function localized at the space-time boundary transforms as the dx^{11} component of a one-form whose other components are identically zero. We will write the gluino condensate at the αth component of the space-time boundary as a four-form $\omega_{(\alpha)}$, in the following coordinate-free way:

$$\omega_{(\alpha)} \equiv \delta(x^{11}) dx^{11} \wedge \langle \overline{\chi^a} \Gamma_{ABC} \chi^a \rangle dx^A \wedge dx^B \wedge dx^C. \tag{3.10}$$

Here $\delta(x^{11})$ is the delta function supported at the αth boundary component and χ's are the corresponding gluinos. For a covariantly constant condensate, $\omega_{(\alpha)}$ is closed and $\mathbf{Z}_2$ invariant, and therefore defines a $\mathbf{Z}_2$-equivariant cohomology class on $\mathbf{X}\times\mathbf{S}^1/\mathbf{Z}_2$. More importantly for our purposes, $\omega_{(\alpha)}$ is closed under the nilpotent operator $d_{11} \equiv dx^{11}\partial_{11}$, which represents the exterior derivative along the 11th dimension, and we denote by $[\omega_{(\alpha)}]$ the corresponding $\mathbf{Z}_2$-equivariant cohomology class in the cohomology defined by d_{11}.

For Eq. (3.5) to have a global solution, the right-hand side of Eq. (3.6) has to be exact with respect to d_{11}, as a $\mathbf{Z}_2$-equivariant form on $\mathbf{X}\times\mathbf{S}^1/\mathbf{Z}_2$. Thus the topological condition that allows us to extend the local solution of the unbroken supersymmetry conditions to a global one is that the $\mathbf{Z}_2$-equivariant d_{11}-cohomology class of $\omega_{(\alpha)}$, summed over all boundary components, vanish:

$$\sum_\alpha [\omega_{(\alpha)}] = 0. \tag{3.11}$$

In general, this condition is violated, and supersymmetry is broken by the global topology of the extra dimension of M theory.

There is, however, one simple way to satisfy the cohomological condition (3.11), which leads to an *a posteriori* plausibility argument indicating that we could have perhaps expected locally unbroken supersymmetry in the presence of a gluino condensate in M theory, with supersymmetry broken only by global topological effects.[8] Set $\widetilde{G}_{IJKL}$ to zero, and consider the case when strong coupling of equal strengths develops in the two gauge groups at the two ends of the world. Now if we put equally strong and opposite gluino condensates at the two boundaries, the topological obstruction (3.11) vanishes, and supersymmetry is unbroken. This can be understood if we start from the limit of weak heterotic coupling, described by ten-dimensional heterotic supergravity. Indeed, it is clear that in the presence of two equally strong gluino condensates that differ only by a minus sign, supersymmetry stays unbroken in ten dimensions. Next, we enlarge the string coupling and go back to the 11-dimensional description. Assuming that the mechanism in which supersymmetry is preserved is local in the 11th dimension, supersymmetry should be locally preserved in the vicinity of each gluino condensate. Now if we change the value of one of the condensates, supersymmetry will be broken, but since it is preserved locally in the neighborhood of each condensate, it can only be broken by effects that involve the global topology of the orbifold dimension. Indeed, it is easy to find the globally defined spinor η'' that represents the unbroken supersymmetry in the background of such equally strong but opposite condensates:

$$\eta'' = \eta_0 - \frac{1}{384\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3} \epsilon(x^{11})(\overline{\chi^a}\Gamma_{ABC}\chi^a)\Gamma^{ABC}\eta_0, \tag{3.12}$$

with η_0 the covariantly constant spinor on $\mathbf{X}$. Clearly, η'' has a jump at both ends of the world, with opposite values corresponding to the strengths of the two gluino condensates.

In the phenomenologically most interesting compactifications—notably, those with the spin connection embedding that breaks one of the E_8's to E_6—one end of the world supports the grand-unified degrees of freedom that are weakly coupled, while the hidden E_8 sector is strongly coupled and should develop a nonzero gluino condensate. With a gluino condensate at only one end of the world, the cohomology condition (3.11) cannot be satisfied, and supersymmetry is broken by the global topology of the extra dimension of M theory, in a mechanism that is reminiscent of the Casimir effect.

[8]This argument was pointed out to me by E. Witten.

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