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PROBLEMS FOR (2,0) COMPACTIFICATIONS

John Ellis, K. Enqvist^{*)}, D.V. Nanopoulos, K. Olive^{**)}, M. Quirós⁺⁾
and
F. Zwirner⁺⁺⁾

CERN - Geneva

ABSTRACT

We expose several theoretical and phenomenological problems which impede the construction of realistic low-energy models based on the superstring using manifolds for which the gauge and spin connections are not identified ($A \neq \omega$), and the world-sheet σ -model has (2,0) supersymmetry. Single world-sheet instantons generate an effective scalar potential which prevents generic manifolds of this type from being solutions to the equations of motion. They also give large supersymmetry-breaking mass parameters $O(m_p)$ to gauge non-singlet fields. The Green-Schwarz modification to the anti-symmetric tensor field strength may not absorb anomalies if $\omega \neq A$. It is difficult to obtain satisfactory mixing among the supersymmetric Weinberg-Salam Higgs fields, which requires an $SU(3) \times SU(2) \times U(1)$ singlet field N with vacuum expectation value $x \equiv \langle 0|N|0 \rangle = O(m_W)$. A Yukawa coupling λN^3 would yield $x = O(m_W)$, but at the expense of cosmologically unacceptable domain walls. If there is no such Yukawa coupling, the natural value for x is $O(m_p)$.

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- *) Permanent address: Research Institute for Theoretical Physics, Univ. of Helsinki, Helsinki, Finland.
**) Physics Dept., University of Minnesota, Minneapolis, Minn., U.S.A.
+) On leave from Instituto de Estructura de la Materia, Madrid, Spain.
++) Also at International School for Advanced Studies, Trieste and INFN, Padova, Italy.

Although many theorists agree that the superstring [1] is the best candidate for a Theory of Everything (TOE), they do not all agree on the best way to make contact with low-energy physics in four dimensions. In particular, several scenarios have been proposed [2-5] for the compactification of the six surplus dimensions. Presumably the effective four-dimensional theory should be supersymmetric, so that one has a chance to solve the hierarchy problem. The requirement of supersymmetry selects compactification manifolds with $SU(3)$ holonomy [2]. Consistency also imposes constraints relating the gauge and spin connections A and ω , to which the most straightforward solution was [2] to identify them: $A = \omega$. The resulting six-dimensional manifolds were the Kähler and Ricci flat Calabi-Yau manifolds [6]. If such a compactification is chosen, the resulting two-dimensional σ -model on the string world-sheet has $(2, 2)$ supersymmetry. It has recently been pointed out that Calabi-Yau manifolds with Ricci flat metrics are not solutions to the equations of motion [7, 8]. However, it is possible to use them as starting points in a perturbative expansion for large compactification scales to obtain manifolds with similar topological properties which are exact solutions. Moreover, it has recently been argued [9a] that in general there exist non-local transformations to non-Ricci flat metrics for the Calabi-Yau manifolds which are solutions to the equations of motion. Hence Calabi-Yau manifolds are theoretically viable compactifications. This argument is supported by the observation [9b] that the central charge of the Virasoro algebra is not modified through 4 loops. However, it is not evident that Calabi-Yau manifolds are the only satisfactory supersymmetric compactifications. One alternative is to consider orbifolds [3], which have isolated singularities, but to retain the identification $A = \omega$. Another alternative is to consider manifolds (or orbifolds) and gauge bundles for which $A \neq \omega$, and the corresponding σ -model^{*)} has $(2, 0)$ supersymmetry [4, 5].

It is the low-energy phenomenology of Calabi-Yau manifolds which has been the most thoroughly studied [10, 11], with some plausible dynamical scenarios being proposed [12]. The four-dimensional gauge group must in this case be some rank 5 or 6 subgroup of E_6 , and the generations of light non-singlet matter particles have the quantum numbers of 27 representations of E_6 [11]. Whilst in no way excluded by experiment, these gauge groups and matter representations are considerably larger than those in the minimal supersymmetric Standard Model, which has a rank 4 gauge group and 15-dimensional matter generations. Therefore, there

*) Note that the arguments of Ref. 9 have been given for $(2, 2)$ models, and not for $(2, 0)$ models.

is considerable interest in whether the additional light four-dimensional fields are inevitable consequences of compactification--and hence a prediction of superstring theory--or whether alternative compactification schemes could avoid introducing extra fields. Indeed, scenarios have been proposed, based on orbifolds [3] and on manifolds with $A \neq \omega$ [4, 5], which could yield four-dimensional gauge groups which are rank 4 subgroups of $SU(5)$ and rank 4 or 5 subgroups of $SO(10)$. The questions then arise whether these compactifications are theoretically consistent and yield phenomenologically viable models. We give some answers to these questions for the manifolds with $A \neq \omega$ which correspond to $(2, 0)$ σ -models.

We show that single world-sheet instantons in $(2, 0)$ σ -models generate a non-trivial effective scalar potential through both D- [13] and F-terms [14]. Therefore generic manifolds with $\omega \neq A$ are not solutions to the equations of motion. Moreover, they lead [13] to unacceptably large supersymmetry-breaking parameters for gauge non-singlet fields. The Green-Schwarz anomaly cancellation mechanism [15] is only proven to work for compactifications in which $A = \omega$. If $A \neq \omega$, the possibility exists [16] of an additional topological difficulty, which must be studied in any specific case. Phenomenological problems centre on the $H\bar{H}$ mixing term which is essential to any supergravity model with realistic electroweak breaking [17]. This can only be provided by an $SU(3)_C \times SU(2)_L \times U(1)_Y$ singlet field N with vacuum expectation value $x \equiv \langle 0|N|0 \rangle = O(m_W)$. Superstring-inspired $SU(5)$ and $SO(10)$ models may contain singlets with the required $H\bar{H}N$ Yukawa coupling [18]. Generically, the N fields have a λN^3 Yukawa coupling as well. If present, it will create three inequivalent vacua of equal energy, and the Universe will have a cosmologically unacceptable [19] structure of domains with excess energy density in domain walls. If the λN^3 coupling is absent, $x = O(\mu_0)$, where μ_0 is the renormalization scale at which the soft supersymmetry-breaking mass-squared M_N^2 goes negative. Since the N field is a gauge singlet, $\mu_0 = O(m_P)$ if the initial value of $|m_N|$ at the Planck scale is much smaller than the supersymmetry-breaking gaugino mass, as we expect [20]. We conclude that it is much more difficult to construct viable models based on $(2, 0)$ compactifications [4, 5] than it was with $(2, 2)$ compactifications [2].

We start by discussing the theoretical problems mentioned above. In the case of Calabi-Yau manifolds with $A = \omega$, the fact that the world-sheet σ -model has $(2, 2)$ supersymmetry means that instantons [21] have equal numbers of left-moving and right-moving zero modes [13, 14]. Therefore they do not induce self-interactions involving only the four-dimensional fields corresponding to moduli of the manifold, since these are all gauge singlets and hence left-movers [13, 14]. On the other hand, instantons in the $(2, 0)$ σ -models associated with $A \neq \omega$ compactifications have only left-moving zero modes, and hence can generate self-interactions among the moduli fields [13, 14]. To the extent that these self-

interactions make a non-trivial contribution to the effective scalar potential, vacua with generic values of the moduli of the compactification manifold will be destabilized. One or more of the moduli may be driven to an extreme value corresponding to some singular orbifold limit of the manifold [3].

In what order do (2, 0) world-sheet instanton effects generate a non-trivial effective scalar potential? Interactions among matter fields in $N = 1$, $d = 4$ supergravity are characterized [22] by a Kähler potential $K = K_D + K_F$, where K_D is a non-holomorphic D-term yielding the Kähler metric $K_{\phi\phi^*} = K_{D\phi\phi^*}$, and K_F is a holomorphic F-term (and its complex conjugate) yielding the superpotential W : $K_F = \ln |W|^2$. All supergravity theories interrelated by the Kähler transformations

$$K_D \rightarrow K_D + \ln |F|^2, \quad W \rightarrow W/F \quad (1)$$

are equivalent, so theories can be divided into equivalence classes of $W = 0$, for which no supersymmetry breaking is possible, and $W \neq 0$. The latter is the class used for realistic model-building [23]. It has been shown that a single world-sheet instanton modifies both K_D [13] and K_F [14]: $\Delta K^{(1)} = \Delta K_D^{(1)} + \Delta K_F^{(1)}$; $\Delta K_D^{(1)}, \Delta K_F^{(1)} \neq 0$. Since the tree-level perturbative Kähler potential $K^{(0)}$ is [10] of the no-scale type [23, 24], $v^{(0)} = 0$. The corresponding one-instanton effect on the effective scalar potential is

$$\Delta V^{(1)} = e^{K^{(0)}} \left[\Delta K_i^{(1)} (K^{(0)-1})^i_j K^{(0)j} - \frac{1}{2} K_i^{(0)} (K^{(0)-1})^i_j \Delta K_k^{(1)j} (K^{(0)-1})^k_l K^{(0)l} \right] + \text{herm. conj.} \quad (2)$$

where $K^{(0)}$ is the tree-level perturbative Kähler potential, plus terms of higher order in the instanton action. In the case of a single modulus T characterizing the overall scale of the compactification manifold, there is a unique [24] form of no-scale Kähler potential giving $V(T) = 0$,

$$K = -3 \ln [f(T) + f^\dagger(T^*)] \quad (3a)$$

which can be cast in the conventional form

$$K = -3 \ln (T + T^*) \quad (3b)$$

by making a holomorphic change of variables: $f(T) \rightarrow T$. For $\Delta V^{(1)}$ to vanish, we would need

$$\Delta K^{(1)} = \frac{g(T) + g^\dagger(T^*)}{T + T^*} \quad (4)$$

which corresponds to just such a holomorphic change of variables. In fact, one finds [13]

$$\Delta K_D^{(1)} = -\frac{1}{2} T^{*2} (T+1) e^{-T} + \text{herm. conj.} \quad (5a)$$

resulting in

$$\Delta V_D^{(1)} = e^{K_D^{(0)}} (T+T^*) \left[\left(1 - T - \frac{3}{2} T^* \right) |T|^2 + T^* \right] e^{-T} + \text{herm. conj.} \quad (5b)$$

whilst [14]

$$\Delta K_F^{(1)} = e^{-K_F^{(0)}} e^{-T} + \text{herm. conj.} \quad (6a)$$

resulting in

$$\Delta V_F^{(1)} = e^{K_D^{(0)}} (T+T^*) e^{-T} + \text{herm. conj.} \quad (6b)$$

Thus the modifications to the Kähler metric (5) and to the superpotential (6) both give a non-trivial effective scalar potential at leading order in the instanton interaction, i.e. in order e^{-T_R} . The modification (6a) to the superpotential W also makes F-term contributions [14] to the effective scalar potential $\propto |\partial W / \partial T|^2$ of order e^{-2T_R} , which are subleading in the limit of the large T_R for which this semiclassical instanton calculation is accurate.

The leading modifications (5b) and (6b) to the effective scalar potential establish that the vacuum is not stable in the limit $T_R \rightarrow \infty$, corresponding to a large size for the manifold of compactification. They suggest that manifolds with generic values of the modulus T are unstable, and not solutions to the equations of motion, but the semiclassical approximation cannot be used to calculate $V(T_R)$ reliably when $T_R = O(1)$. It is known that orbifolds [3] are solutions to the equations of motion, and may correspond to singular limits of manifolds when one or more of the moduli vanish, but the above shortcoming of the semiclassical approximation prevents one from being sure that $(2, 0)$ compactification inevitably collapse to orbifolds. The most one can say is that $(2, 0)$ compactifications are generically unstable.

Another major difficulty with the modifications (5, 6) to the effective scalar potential arises when gauge non-singlet fields ϕ^i are included. As was pointed out in Ref. 13, such modifications $\Delta K^{(1)}$ generate supersymmetry-breaking scalar masses and trilinear scalar couplings with dimensional coefficients which are $O(m_P)$. If present, these would prevent the establishment of a gauge hierarchy

$m_W \ll m_P$. Model-builders using (2, 0) compactifications need to assume that these can be avoided somehow, as we do for the sake of argument in the rest of this paper. In sharp contrast, instanton effects in (2, 2) models generate a $\Delta V(\phi)$ which vanishes for all T when the conventional F - and D -terms vanish, thus maintaining a flat direction with no cosmological constant [13]. In particular, $\Delta V(\phi)$ does not contain any soft supersymmetry-breaking terms for the non-singlet fields, and thus is consistent with the generation of a gauge hierarchy $m_W \ll m_P$ through higher-order radiative corrections in some variant of a no-scale scenario [12].

A third possible difficulty for (2, 0) compactifications goes back to the anomaly cancellation mechanism of Green and Schwarz [15]. They suggested modifying the field strength H of the antisymmetric tensor B to include both gauge and Lorentz Chern-Simons terms:

$$H = -dB + \omega_G^3 - \omega_L^3 \quad (7)$$

This procedure is consistent if $(\omega_G^3 - \omega_L^3 - H)$ is in the trivial de Rham cohomology class of $H^3(P)$, where $P(M, G)$ is the principal fibre bundle on the base space M with fibre the Lie group G . The presence of ω_L^3 in Eq. (7) is necessary, since $(\omega_G^3 - H)$ is related to the non-trivial generator of $H^3(G)$, and hence is a non-trivial element of $H^3(P)$ [16]. If the gauge and spin connections on M are identified, then $(\omega_G^3 - \omega_L^3 - H)$ is evidently a trivial element of $H^3(P)$, and all is well [16]. However, if the gauge and spin connections are not equal, the possibility that $(\omega_G^3 - \omega_L^3 - H)$ is non-trivial resurfaces [16] and the procedure (7) might not be consistent. This must presumably be studied in each case, and may be a significant constraint on compactifications with unequal gauge and spin connections corresponding to (2, 0) σ -models.

We now turn to some phenomenological problems related to electroweak symmetry breaking. So far, the only scenario for triggering this is dimensional transmutation [12], with one supersymmetry-breaking Higgs squared mass being driven negative by Yukawa couplings [17]. The minimal supersymmetric Standard Model contains two isodoublet Higgs fields H and $\bar{H}$, which must be coupled in a superpotential term $\epsilon H \bar{H}$ if both $v \equiv \langle 0|H|0 \rangle$ and $\bar{v} \equiv \langle 0|\bar{H}|0 \rangle$ are to be non-zero. This also avoids the appearance of an unacceptable axion. To have both $v, \bar{v} = O(m_W)$, the mixing coefficient ϵ must also be $O(m_W)$. Phenomenological superstring models do not contain bilinear superpotential terms, so ϵ must be obtained from the vacuum expectation value of some $SU(3)_C \times SU(2)_L \times U(1)_Y$ singlet field N :

$$W \ni \lambda_M N H \bar{H}, \quad \langle 0|N|0 \rangle = \chi \Rightarrow \epsilon = \lambda_M \chi \quad (8)$$

and to get $\epsilon = 0(m_W)$, also x must be $0(m_W)$. In $(2, 2)$ compactifications, matter fields occupy 27 representations of E_6 which contain a neutral field N able to do the job (8) [12]. In $(2, 0)$ compactifications, the matter fields may contain a singlet of $SU(5)$ or $SO(10)$ (but not an adjoint 24 or 45), which would be a candidate for the job (8) [4, 18]. However, it is not clear that an $SU(5)$ or $SO(10)$ singlet field N can in fact do the job (8), as we now discuss.

The original 10-dimensional E_6 gauge interactions will in general generate a trilinear self-coupling of the N field:

$$W \supset \lambda N^3 \quad (9)$$

The corresponding effective low-energy renormalizable potential for the N field is

$$V(N) = m_N^2 |N|^2 + (Am_N \lambda N^3 + h.c.) + |\lambda N^2|^2 \quad (10)$$

where m_N is a soft supersymmetry-breaking scalar mass of order m_W and A is a numerical coefficient of order 1. Clearly the potential (10) has a Z_3 discrete symmetry which must be spontaneously broken in the vacuum, threatening a domain wall problem. For simplicity, we take λ to be real and write $N = x e^{i\phi}$ so that V (10) takes the form:

$$V(x, \phi) = m_N^2 x^2 + 2Am_N \lambda \cos 3\phi x^3 + \lambda^2 x^4 \quad (11)$$

We can now distinguish two essentially different cases: $\lambda \neq 0$ as generally expected, and $\lambda = 0$.

When $\lambda \neq 0$ and $|A| > 1$, the potential (11) exhibits three inequivalent global minima of equal vacuum energy^{*}, with

$$\frac{x}{m_N} = \frac{3|A| + \sqrt{9A^2 - 8}}{4\lambda}; \quad \phi = \begin{cases} 0, \pm \frac{2\pi}{3} & (A\lambda < 0) \\ \pi, \pm \frac{\pi}{3} & (A\lambda > 0) \end{cases} \quad (12)$$

^{*}) Note that the $H\bar{H}N$ superpotential term (8) does not modify this discussion. The Z_3 symmetry of the vacua (12) can be extended to allow for $\langle 0|H, \bar{H}|0\rangle \neq 0$ by including corresponding discrete phase transformations on $H, \bar{H}$. These do not modify electroweak physics because they are included in the electroweak group. Similarly, the Z_3 vacuum problem does not arise in an E_6 model in which N carries an extra $U(1)_E$ hypercharge, and there is no N^3 term in the superpotential.

If $|A| < 1$, the only global minimum of (11) is at $x = 0$, which is phenomenologically unacceptable. As the Universe cools below temperatures $T = O(x) = O(m_W)$, domains of order of the horizon size which sit in the different vacua (12) form, with domain walls between them carrying a surface energy density

$$\sigma \sim \lambda x^3 \quad (13)$$

Domains will coalesce as the Universe expands, but there will always be $O(1)$ domain wall per horizon. If we demand that the energy E_D in one domain wall today be less than the present matter density ρ_0 . i.e.

$$\lambda x^3 / t_0 \lesssim \rho_0 \quad (14)$$

where $t_0 \sim 5 \times 10^{17}$ s is the present age of the Universe, we find an upper bound [19]

$$x \lesssim \frac{40 \text{ MeV}}{\lambda^{1/3}} \quad (15)$$

But Eq. (12) tells us that

$$x \sim m_W / \lambda \quad (16)$$

which is grossly incompatible with the requirement (15). This conclusion can be strengthened by demanding that the domain wall does not make an unacceptable perturbation in the microwave background radiation: $E_D \lesssim (\delta\rho/\rho)\rho_0 \sim 10^{-4} \rho_0$, in which case one finds [19]

$$x \lesssim \frac{2.7 \text{ MeV}}{\lambda^{1/3}} \quad (17)$$

We conclude that the superpotential (9) gives unacceptable domain walls.

One might hope to evade this conclusion by modifying the potential (10) while keeping $\lambda \neq 0$ in the superpotential (9). The strategy would be to 'tilt' the potential (10) so that the inequivalent minima (12) have different energies: $\Delta V = \epsilon \lambda^2 x^4$, $1 > \epsilon > 0$, and one of the minima (10) is dynamically favoured [19]. This could be done by adding non-renormalizable terms to the superpotential W , such as

$$W \supset \lambda_n N^{3n} / m_P^n : n > 0 \quad (18)$$

One way to depopulate the undesired vacua is by tunnelling to the lowest minimum. The tunnelling action at the optimum epoch when $T = O(\lambda x)$ is given in the thin-wall approximation by

$$A \sim \frac{16\pi}{3\epsilon^2} \quad (19)$$

which is acceptably small ($A \lesssim 180$) if $\epsilon > 0.3$. This is clearly not possible for any modified superpotential (18) with $n > 0$, since $x = O(m_W) \ll m_P$. Another way to depopulate the undesired minimum is when the total excess energy in a higher domain begins to exceed the energy in the domain wall surrounding it, so that it becomes energetically favourable to expand the wanted domains while shrinking the unwanted domains, and there is no barrier to prevent this [19]. This occurs at a time $t_* \sim \sigma/\Delta V$. To be useful, this mechanism must operate before the Universe's energy density becomes dominated by domain walls, which occurs at $t_c \sim \sigma^{-1} m_P^2$. Requiring $t_* < t_c$ we find

$$\Delta V \gtrsim \sigma^2 \approx \lambda^2 x^6 / m_P^2 \quad (20)$$

This will occur if $n = 1$ is present, and possibly even if $n = 1$ is absent but $n = 2$ is present, among the non-renormalizable superpotential terms (18). This provides a possible avenue for the construction of a viable $(2, 0)$ model, although one should first address oneself to the question where such a non-renormalizable term might come from. It cannot be generated by σ -model perturbation theory [4], nor by instantons in $(2, 0)$ models [13, 14]. It would have to be generated by integrating out massive string modes, a scenario whose feasibility would have to be demonstrated, not merely postulated.

Finally we turn to the case where $\lambda = 0$ in equations (7) to (9). In this case, the only renormalizable term in the potential for N is

$$V(N) = m_N^2 |N|^2 \quad (21)$$

The only way to get $x = \langle 0|N|0 \rangle \neq 0$ is by dimensional transmutation [12] with the renormalization group equations driving m_N^2 negative at some renormalization scale μ_0 , in which case $x = O(\mu_0)$. Since N is a gauge singlet field, the only contributions to $\mu dm_N^2/d\mu$ come from Yukawa interactions such as λ_M (8), and $\mu dm_N^2/d\mu < 0$ [12]. If the leading source of supersymmetry breaking at the Planck scale is the gaugino mass $m_{1/2}$, as we believe [20], then $\mu_0 = O(m_P)$ and hence

$$x = O(m_P) \quad (22)$$

which is unacceptable.

The phenomenological problems we have presented above do not arise in 4-dimensional models based on rank 5 or 6 subgroups of E_6 , as come from (2, 2) compactifications. In these models, the Z_3 supersymmetry of the potential is embedded in the gauge group, the three previously inequivalent vacua are linked by gauge transformations, and there are no domains or domain walls. Moreover, the gauge interactions of the N field make a positive contribution to $\mu d\mu_N^2/d\mu$, which ensures that μ_0 is hierarchically smaller than m_P and $x = O(m_W)$ is natural [12].

Although we have concentrated in this paper on manifolds corresponding to (2, 0) σ -models [4, 5] which lead to the 4-dimensional gauge groups SU(5) and SO(10), analogous phenomenological problems with Higgs mixing may also arise in orbifold compactifications [3] which lead to SU(5) or SO(10). Maybe we phenomenologists will have to stick to E_6 .

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