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WORLD SHEET INSTANTON EFFECTS ON NO-SCALE STRUCTURE

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ABSTRACT

We study the effects of world sheet instantons on the effective low-energy four-dimensional supergravity theory obtained by compactification of the heterotic superstring on a Calabi-Yau manifold K_6 . World sheet instantons do not modify the Kähler metric for the gauge-singlet S field, nor do they modify the kinetic terms for the gauge fields. They modify the Kähler metric for the T field if $\dim H_1(\text{End } T_6) \neq 0$, where T_6 is the tangent bundle on K_6 , whereas they modify the Kähler metric for the gauge non-singlet fields ϕ^i if $\dim H_1(\text{End } T_6) = 0$. If present, the modifications to the Kähler metric of the T field would create a non-trivial effective scalar potential destroying its flat no-scale structure, and would in principle enable $\langle 0|T|0 \rangle$ to be dynamically determined to be of $O(1)$, though we are not able within our approximations to calculate its value reliably. The self-interactions of the gauge non-singlet fields ϕ^i maintain their no-scale structure.

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The earliest approach taken to deduce the form of the effective low-energy four-dimensional field theory obtained from the superstring¹⁾ was first to take the zero-slope limit in which the string excitations were integrated out to leave a 10-dimensional supergravity theory²⁾, and then to compactify this field theory on a Calabi-Yau manifold or reasonable facsimile thereof³⁾.

The resulting low-energy four-dimensional theory took the form of a no-scale supergravity⁴⁾ theory with a Kähler potential⁵⁾

$$\mathcal{G} = - \ln (S + S^\dagger) + G + \rho_n |P(S, \phi^i)|^2 \quad (1a)$$

where⁴⁾

$$G = -3 \rho_n (\tau + \tau^\dagger - 2 \phi^i \phi_i^\dagger) \quad (1b)$$

and a gauge field kinetic term⁵⁾

$$\int d^2\theta \, S W_\alpha W_\alpha + (\text{h.c.}) \quad (2)$$

where S and T are gauge-singlet fields, the ϕ^i are gauge non-singlets, P is the superpotential and W_α the chiral gauge superfield. One might expect the no-scale structure (1) and (2) to be modified by many different types of corrections, and attempts^{6),7)} have been made to study some of them. For example, it has been argued⁷⁾ that the local Green-Schwarz term⁸⁾ introduced to cancel the anomaly generated by one loop diagrams in the 10-dimensional supergravity theory does not alter⁷⁾ the gauge field kinetic term (2). In principle, one could look at other loop effects in the supergravity theory in 10 dimensions. However, the relevance of the 10-dimensional supergravity theory²⁾ as an intermediate step has been questioned⁹⁾, and one would prefer to go directly from the superstring in 10 dimensions to the field theory in four dimensions. Arguments have been given^{3),10)} that Calabi-Yau manifolds K_6 are still of interest for compactification, and that the numbers of massless gauge-singlet and gauge-non-singlet fields are determined by topological invariants such as the Betti numbers^{3),11)}. It is also known¹¹⁾ that higher order terms in the topological expansion for superstring amplitudes do not renormalize the superpotential P in Eq. (1). A general framework has been set up¹²⁾ for determining what aspects of the $[SU(1,1)/U(1)]_S \times [SU(n+1,1)/SU(n+1) \times U(1)]_{(\phi^i, T)}$ structure of the Kähler potential (1) may be preserved or destroyed by non-perturbative world sheet instanton effects¹³⁾ at any given order in the topological expansion for the superstring. In particular, we

have argued¹²⁾ that world sheet instantons¹³⁾, which are the only known non-perturbative effects at the tree level of the superstring, would in general break the $SU(1,1)/U(1)$ subsymmetry associated with the T field in (1b).

In this paper the general framework of Ref. 12) is used to characterize the effects of world sheet instantons on the Kähler potential (1) and on the gauge field kinetic term (2). We find no modification of the Kähler metric for the S field, but find a modification of the Kähler metric for the T field if and only if $\dim H_1(\text{End } T_6) \neq 0$, where T_6 is the tangent bundle on K_6 . The modification of the Kähler metric for the gauge non-singlet fields ϕ^i found in Ref. 13) is present if and only if $\dim H_1(\text{End } T_6) = 0$. We also argue that world sheet instantons do not modify the simple form of the gauge kinetic term (2). The new interactions do not modify the effective scalar potential unless supersymmetry is broken, in which case the new T interactions (if present) provide a non-flat potential $\propto m_{3/2}^2$. Thus the no-scale structure for the T field is indeed destroyed if $\dim H_1(\text{End } T_6) \neq 0$, whereas the ϕ^i interactions by themselves would not break the no-scale structure if $\dim H_1(\text{End } T_6) = 0$. A full determination of the T potential could enable one to fix the value of the T field, but our calculation is not valid for $T = 0(1)$ where the physical value presumably lies.

We consider world sheet instanton effects¹³⁾ in the $E_8 \times E_8$ ' heterotic string¹⁴⁾ compactified on $M^4 \times K_6$, where the six-dimensional manifold K_6 has a non-zero second Betti number b_2 . In this case the antisymmetric tensor field $B_{ij}(x, \xi)$: $i, j = 1, 2, \dots, 6$, $(x, \xi) \in M^4 \times K_6$ can be written in terms of the b_2 harmonic forms $b_{ij}^a(\xi)$: $a = 1, 2, \dots, b_2$, the basis of $H_2(K_6)$, as follows¹⁵⁾:

$$B_{ij}(x, \xi) = \sum_{a=1}^{b_2} \phi^a(x) b_{ij}^a(\xi) \quad (3)$$

When $b_2(K_6) = 1$, the field $\phi(x)$ in Eq. (3) is just the imaginary part T_I of the T superfield. World sheet instantons are topologically non-trivial string configurations whose integrated Wess-Zumino term does not vanish¹³⁾:

$$\nu(X) \equiv \int_{\Sigma} \epsilon^{\mu\nu} b_{ij}^a(X) \partial_{\mu} X^i \partial_{\nu} X^j \neq 0 \quad (4)$$

This "topological charge" will be non-zero for some topologically non-trivial string configuration X whenever the harmonic form b_{ij}^a is not exact. It is

clear from the definition (4) of the world sheet instanton that the configuration is completely contained in K_6 , and will appear already at the tree level in the string loop expansion if $\pi_2(K_6) \neq 0$.

This kind of instanton¹³⁾ induces effective interactions on the world sheet in the same way as ordinary instantons in M^4 . Our interactions take the form

$$e^{i\phi(x)} e^{-S_{cl}} \prod_{i=1}^k \Psi_i^0(\sigma) \quad (5)$$

where we have specialized to the case $b_2 = 1$, and S_{cl} is the classical action for the string instanton configuration. The fermions ϕ_i can be visualized as two-dimensional Ramond-Neveu-Schwarz fermions¹⁶⁾ defined on the string world sheet. Equation (5) contains the product of the wave functions ϕ_i^0 of the k fermionic zero modes. In the case of non-holomorphic instantons, the number of fermionic zero modes can be obtained¹³⁾ by acting on the string configuration with the supersymmetric charges of the non-linear σ -model defined on the string world sheet. The existence of these world sheet zero modes implies that this kind of instanton can only give expectation values for products of vertices with enough fermions to eat up all the world sheet zero modes. These non-zero expectation values will be our starting points for deriving the four-dimensional effective couplings induced by the instantons.

The counting of fermionic zero modes for non-holomorphic instantons reduces to counting the numbers of left and right supersymmetries of the world sheet σ -model. In the heterotic string before compactification it is an $N = \frac{1}{2}$ σ -model¹⁴⁾ with just one right supersymmetry. However, if we compactify our string theory on $M^4 \times K_6$, the string σ -model must have two supersymmetries in order to retain at least one supersymmetry in four dimensions after compactification^{3),17)}. This means that the string σ -model for the heterotic string must be of the (2,0) type with two right supersymmetries^{3),17)}. Using the fermionic representation, the corresponding (2,0) σ -model in a gauge background fibre bundle on K_6 is¹⁷⁾

$$S = \int d^2\tau d^2\theta \left\{ [G_{ij} + B_{ij}] DX^i \frac{\partial}{\partial \tau} X^j - F_{AB}(X) \lambda^A (D\lambda^B + A_{ic}^B (DX^i) \lambda^c) \right\}_{(6)}$$

where we use Neveu-Ramond-Schwarz fermions¹⁶⁾ and the λ^A : $A = 1, 2, \dots, 32$ are the superfields associated with the two sets of left-handed fermions used to fermionize the internal set of bosonic string co-ordinates. The superfields

X^i have as their fermionic components the right-moving Neveu-Ramond-Schwarz fermions¹⁶⁾. The gauge connection A_C^B and the field strength F_{AB} in Eq. (6) are specified by the gauge fibre bundle on K_6 . The (2,0) σ -model defined by (6) is free of anomalies if we allow^{17),18)} the antisymmetric tensor field B_{ij} to transform à la Green and Schwarz⁸⁾:

$$\delta B = \text{tr} (A \wedge \Lambda) - \text{tr} (\omega \wedge \theta) \quad (7)$$

where ω is the spin connection on K_6 . As is well known^{3),19)}, Eq. (7) imposes the following constraint for anomaly freedom on K_6 :

$$\text{tr} (F \wedge F) = \text{tr} (R \wedge R) \quad (8)$$

to which the most natural solution is

$$A = \omega \quad (9)$$

corresponding to the embedding of the $SU(3)$ spin connection into E_8 , and the breaking of E_8 to $SU(3) \times E_6$. If we choose this particular solution to the anomaly freedom constraint (8), then the σ -model (6) becomes left-right symmetric with (2,2) supersymmetry^{14),17)}. For non-holomorphic instantons, Eq. (9) implies two left and two right fermionic zero modes.

For holomorphic instantons defined by the self-duality condition¹³⁾

$$S(X_{\text{inst}}) = \nu(X_{\text{inst}}) \quad (10)$$

we get, by acting with the supersymmetry charges of the σ -model, one left and one right zero mode. If the instanton appears at the tree level, i.e., on S_2 because $\pi_2(K_6) \neq 0$, we must multiply by two the number of fermionic zero modes as a consequence of the superconformal invariance of the model.

Next we look at the vertices for the low mass superfields of interest. The S field is dismissed briefly: its vertex only contains indices tangent to M^4 , not to K_6 . Therefore, it cannot soak up any of the string σ -model zero modes discussed above, and world sheet instantons generate no new interactions involving the S field. The vertices for the superfield T , whose real part T_R

is the dilaton controlling the size of the internal manifold K_6 , take the forms²⁰⁾

$$V(T_R) = e^{iK \cdot X} G_{ij}(X) \partial^\mu X^i \partial_\mu X^j + \text{susy completion} \quad (11a)$$

$$V(T_I) = e^{iK \cdot X} b_{ij}^0(X) \partial_\mu X^i \partial_\nu X^j \epsilon^{\mu\nu} + \text{susy completion} \quad (11b)$$

where $i = 1, 2, \dots, 6$. In the case of the heterotic string, the susy completions of (11a) and (11b) are

$$e^{iK \cdot X} G_{ij}(X) \partial X^i \frac{1}{2} K_A \frac{1}{8} \bar{\chi} \gamma^{Aj} \chi \quad (12a)$$

$$e^{iK \cdot X} b_{ij}^0(X) \partial X^i \frac{1}{2} K_A \frac{1}{8} \bar{\chi} \gamma^{Aj} \chi \quad (12b)$$

where the χ are 10-dimensional right-moving fermions. It is important to emphasize that the vertices for all gauge singlet superfields only contain right-moving fermions. Left-moving fermions could only appear using the fermionic representation of the internal momenta $P^I = \partial X^I / \partial \tau^-$: $I = 1, 2, \dots, 16$ on the compactified 16-dimensional torus¹⁴⁾. The vertex for a gauge non-singlet superfield ϕ^i was given in Ref. 13): we do not give it again here, but recall that it contains one left-moving and one right-moving fermion with indices tangent to K_6 . Finally, we dismiss the gauge superfield W_α for the same reasons as the S superfield: its vertex only contains fermions with indices tangent to M^4 , and hence world sheet instantons do not induce any new interactions involving W .

We are now in a position to study the new interactions of T fields induced by world sheet instantons. In fact, we see immediately that when we choose the simple solution (9) to the anomaly cancellation condition (8), i.e., when we break E_8 to $SU(3) \times E_6$, we cannot get a non-zero expectation value for any product of vertices for the dilaton superfield T . This is because these vertices do not contain the left-moving fermions which we need to soak up the whole set of fermionic zero modes.

However, a solution to the anomaly cancellation condition (8) which has $A \neq \omega$ can be obtained¹¹⁾ if $\dim H_1(\text{End } T_6) \neq 0$, where T_6 is the $SU(3)$ gauge fibre bundle defined on K_6 by the embedding (9). In this case, we can use different gauge connections in the moduli of $A = \omega$ which have the same values

for all topological invariants, including in particular the value of the Pontryagin number $\text{tr } F \wedge F$ which appears in (8). Using these deformations we can break the left-right symmetry of our σ -model. Then we will have a (2,0) σ -model, which means that non-holomorphic instantons at least will only have right-moving fermionic zero modes. Therefore, if E_8 is not broken to $SU(3) \times E_6$ we can get non-zero expectation values for products of vertices for gauge-singlet fields, and hence new effective couplings for these four-dimensional fields.

Let us now consider the minimum non-zero expectation value we can get for T fields in the presence of a non-holomorphic instanton for a (2,0) σ -model, i.e., when $A \neq \omega$. Remembering that the instanton lives in K_6 , and using Neveu-Ramond-Schwarz fermions¹⁶⁾ for the vertices (11) and (12), we see that each T vertex can only soak up one right-handed zero mode if $b_2(K_6) = 1$ and we identify the harmonic form b_{ij}^0 with the Kähler form which is covariantly constant. At the tree level, i.e., on S^2 , we have four right-moving zero modes on a non-holomorphic instanton, because of superconformal symmetry. Hence the effective four-dimensional coupling that we can get is

$$\Delta \mathcal{L}_{\text{eff}}^{(\tau)} \sim e^{-S_a} (t \gamma_\lambda t t^\dagger \gamma^\lambda t^\dagger) e^{i T_I} \quad (13)$$

where $t_R = T_R - |\phi|^2$ is related to the size of the internal manifold, $t_I = T_I$ and we have normalized $v(X_{\text{inst}}) = 1$. Note that, as observed in Ref. 13), the lowest order effective coupling induced by world sheet instantons must vanish if any external boson has zero momentum. This can be seen by direct inspection of the vertices (11) and (12): when $k = 0$ they reduce to parts of the string action. But the new amplitude with one less vertex is to be computed for a Lagrangian with the same amount of supersymmetry on the world sheet, and hence the same number of fermionic zero modes. This amplitude is clearly zero because now we do not have enough vertices to soak up all the world sheet fermionic zero modes.

As for other effective couplings induced by instantons, Wen and Witten have already observed¹³⁾ that a quartic interaction for the gauge non-singlet fields ϕ^i can be induced:

$$\Delta \mathcal{L}_{\text{eff}}^{(\phi)} \sim f_e f^{\dagger e} e^{-S_a} e^{i T_I} : f \equiv d_{emn} \phi^e \phi^m \phi^n \quad (14)$$

It is clear from the above counting of fermionic zero modes that such a term is only generated if $\dim H_1(\text{End } T_6) = 0$ so that one has a (2,2) σ -model on the

world sheet. Moreover, it is clear that no other quartic interactions mixing T and ϕ^i are induced. Also, as remarked above, there are no new couplings involving the S and W superfields because their vertices involve world sheet fermions with M^4 indices.

We now study the phenomenological ramifications of the new couplings (13) and (14). They can both be expressed in terms of modifications ΔG to the original no-scale Kähler potential $\mathcal{G}$ (1). In the absence of a calculation of higher order couplings induced by a single instanton and of multi-instanton effects, we treat ΔG as a perturbation on $\mathcal{G}$. The resulting change in the leading order potential

$$\begin{aligned} V_0 &= e^{\mathcal{G}} \left(\mathcal{G}^i (\mathcal{G}^{-1})^{\bar{j}} \mathcal{G}_{\bar{j}} - 3 \right) + \frac{1}{2} \mathcal{D}^\alpha \mathcal{D}^\alpha = \\ &= \frac{1}{2s_R} e^{\mathcal{G}} \left\{ |\Gamma|^2 + \frac{1}{6} e^{-\mathcal{G}/3} |\mathcal{P}_e|^2 \right\} + \frac{1}{2} \mathcal{D}^\alpha \mathcal{D}^\alpha \end{aligned} \quad (15)$$

where

$$\Gamma \equiv \mathcal{P}(s, \phi^i) - 2s_R \mathcal{P}_s \quad (16)$$

is

$$\begin{aligned} \Delta V &= V_0 \Delta G + \\ &+ e^{\mathcal{G}} \left(\Delta \mathcal{G}^i (\mathcal{G}^{-1})^{\bar{j}} \mathcal{G}_{\bar{j}} - \mathcal{G}^i (\mathcal{G}^{-1})^{\bar{j}} (\Delta \mathcal{G})_{\bar{j}} (\mathcal{G}^{-1})^{\bar{k}} \mathcal{G}_{\bar{k}} + \mathcal{G}^i (\mathcal{G}^{-1})^{\bar{j}} \Delta \mathcal{G}_{\bar{j}} \right) \end{aligned} \quad (17)$$

The leading order potential V_0 (15) is reminiscent⁴⁾ of global supersymmetry, and has a minimum with zero cosmological constant at

$$\Gamma = \mathcal{P}_e = \mathcal{D}^\alpha = 0 \quad (18)$$

Equation (18) defines a flat direction, since $V_0 = 0$ while the vacuum expectation value of T is undetermined. Although local supersymmetry is spontaneously broken in the minimum (18), the scalars ϕ^i do not acquire any soft supersymmetry breaking mass⁴⁾ from (15), nor are there trilinear or quadrilinear interaction terms for the ϕ^i which break supersymmetry⁴⁾. One can wonder whether these properties are preserved by ΔV , since ΔG contains non-trivial couplings between the T and ϕ^i fields. These do not contradict the previously quoted observation that the interactions induced by instantons should vanish if any external boson has zero momentum, since the effective potential couplings in (17) vanish in the supersymmetric limit: $m_{3/2} =$

$e^{g/2} \rightarrow 0$. It was already remarked in Ref. 13) that a mass $m_{T_I} = (\partial^2 V / \partial T_I^2)^{1/2}$ can be generated if $m_{3/2} = e^{g/2} \neq 0$.

We consider first the form of ΔV generated if $\dim H_1(\text{End } T_6) \neq 0$ by an addition to the Kähler potential G of the form

$$\Delta G = \epsilon K(T, T^\dagger, |\phi|^2) \quad , \quad \epsilon \propto e^{-S_Q} \quad (19)$$

which is

$$\begin{aligned} \Delta V = & \frac{1}{2} V_0 K + \frac{1}{2S_R} e^{\frac{2}{3}G} e^{P+P^\dagger} \left\{ -K_T + \frac{1}{3} K_{T^\dagger} \phi^e P_e + \right. \\ & + \frac{1}{6} K^e P_e + e^{-G/3} \left[K_{T T^\dagger} \left(-\frac{1}{2} - \frac{1}{18} (\phi^e P_e)^2 + \right. \right. \\ & \left. \left. + \frac{1}{3} \phi^e P_e \right) + \frac{1}{6} K_T^e P_e \left(2 - \frac{1}{3} \phi_m^\dagger P^m \right) - \frac{1}{36} K_m^e P^m P_e \right] \right\} + \text{h.c.} \end{aligned} \quad (20)$$

Clearly ΔV (20) will in general have a non-trivial dependence on T when $|\phi| = 0$, so the no-scale flatness²¹⁾ of the effective potential in the T direction is broken. This gives hope that a dynamically preferred value of T could be found by minimizing ΔV . Such a hope is intuitively reasonable, since the value of T_R determines the size of the compactified dimensions, and the action of the world sheet instanton is bounded by the area of its world surface, so one might expect instanton effects to prefer a certain size of the compactified dimensions and hence a value of T_R . However, a general function $K(T, T^\dagger, |\phi|^2)$ in (19) will give a minimum of ΔV with non-zero cosmological constant, soft supersymmetry breaking terms [cf. the third, fourth, seventh and eighth terms in (20)] and perhaps even supersymmetry breaking quadrilinear terms [cf. the ninth term in (20)].

The interaction (13) corresponds to the particular case

$$K(t, t^\dagger) = (16 - 4 t t^\dagger) \cos T_I + 16 T_I \sin T_I + K(T) + K^\dagger(T^\dagger) \quad (21)$$

when instantons and anti-instantons are added together, where $k(T)$ is an analytic function of T . If $k(T) = 0$, and assuming²²⁾

$$P(S, \phi^i) = c + h \exp\left(-\frac{8\pi^2}{C_2(G)} S\right) + \hat{P}(\phi^i) \quad (22)$$

at the global minimum of V_0 given by the conditions (18) we have

$$V = \epsilon \frac{\pi}{\alpha_{\text{GUT}}} \left(\frac{2\pi h}{C_2(G)} \right)^2 \exp \left(- \frac{4\pi}{\alpha_{\text{GUT}} C_2(G)} \right) \mathcal{V}(T_R, T_I) \quad (23)$$

where

$$\mathcal{V}(T_R, T_I) = \frac{4 - T_R^2 - T_I^2}{T_R} \cos T_I \quad (24)$$

The reduced potential (24) has no minima or maxima, only saddle points, but this need not concern us. It is easy to find an analytic function $k(T) \neq 0$ which gives V a local minimum at $T_R = 0(1) \neq 0$, $T_I = 0$ with no cosmological constant, but we do not attach much significance to this. We expect the ensemble of non-perturbative effects to fix a dynamical value of $T = 0(1)$, and the lowest order instanton interaction (13), (19), (21) and (24) certainly tells us that $T = 0$ is unstable, just as previous perturbative calculations^{23),24)} in the effective low energy theory told us that $1/T = 0$ is unstable²⁴⁾. However, those perturbative calculations were not reliable²⁴⁾ when $T = 0(1)$, and neither does our interaction (13), (19), (21), (24) give a reliable approximation to all the non-perturbative effects when $T = 0(1)$. We can expect the tree level world sheet instanton to yield interactions of higher order in T , instantons appearing at higher order in the topological expansion of the string world sheet will in general give different forms of T interactions, as will multi-instanton configurations, etc. Thus, although we have shown that world sheet instanton effects break the no-scale structure of the naive Kähler potential (1b) for the T field, and can see how a value of $T = 0(1)$ can be dynamically determined, we are not able to calculate its value reliably.

Finally we turn to the possible interactions (14) for the ϕ^i fields, if $\dim H_1(\text{End} T_6) = 0$, which correspond to

$$\Delta G^{(\phi)} = \epsilon f_{ef}^2 \cos T_I \quad (25)$$

yielding

$$\begin{aligned}
\Delta V^{(\phi)} = & \frac{\epsilon}{12} |f_e|^2 V_0 \cos T_I + \\
& + \frac{\epsilon}{24 S_R} e^{\frac{1}{3} G + \hat{P} + \hat{P}^+} \left\{ \cos T_I \left[\frac{1}{3} e^{G/3} f^m f_n \hat{P}_e + \right. \right. \\
& + \frac{1}{4} |f_e|^2 \left(1 - \frac{2}{3} \hat{P}_e \phi^e + \frac{1}{9} |\phi^e \hat{P}_e|^2 \right) - \frac{1}{72} |\hat{P}^n f_{ne}|^2 \Big] + \\
& + \frac{i}{3} \sin T_I \left[e^{G/3} |f_e|^2 \phi^n \hat{P}_n + \frac{1}{2} f^{en} \hat{P}_e f_n + \right. \\
& \left. \left. + \frac{1}{6} \hat{P}^m f_{mn} f^n \phi^e \hat{P}_e \right] \right\} + \text{h.c.} \tag{26}
\end{aligned}$$

The expression (26) includes non-trivial non-renormalizable interactions between T and the matter fields ϕ^i . However, this $\Delta V^{(\phi)}$ vanishes for all T at the global minimum (18) of V_0 , so that the instanton effects (14) maintain a flat no-scale direction in the potential with vanishing cosmological constant.

Moreover, $\Delta V^{(\phi)}$ (26) does not include any soft supersymmetry breaking terms for the matter fields. Thus the no-scale properties⁴⁾ of the effective low energy theory (1) remain unchanged in the presence of the addition (25) to the Kähler potential implied by the ϕ^i interaction induced by world sheet instantons.

In conclusion, we have seen that world sheet instantons do not change the S - or W_α -dependence of the effective four-dimensional supergravity theory (1) and (2) abstracted naively from the superstring. However, world sheet instantons do in general induce non-trivial interactions for the T and ϕ^i fields which depend on the nature of H_1 (End T_6). These interactions may enable the value of the T field to be determined dynamically, but need not modify the effective low energy interactions of the ϕ^i fields which were previously inferred from the no-scale Kähler potential (1). Clearly world sheet instantons play a crucial rôle in the dynamics of superstring theories.

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