ABSTRACT

We study cosmological implications of the duality $\left(PSL(2, \mathbb{Z})\right)$ invariant potential for the compactification radius $T$, arising in a class of superstring vacua. We show that in spite of having only one minimum in the fundamental domain of the $T$ field there are two types of non-supersymmetric domain walls: one is associated with the discrete Peccei-Quinn symmetry $T \rightarrow T + i$, analogous to the axionic domain wall, and another one associated with the noncompact symmetry $T \rightarrow 1/T$, analogous to the $Z_2$ domain walls. The first one is bound by stringy cosmic strings. The scale of such domain walls is governed by the scale of gaugino condensation ($\mathcal{O}(10^{16} \text{ GeV})$ in the case of hidden $E_8$ gauge group), while the separation between minima is of order $M_{pl}$. We discuss the formation of walls and their cosmological implications: the walls must be gotten rid of, either by chopping by stringy cosmic strings and/or inflation. Since there is no usual Kibble mechanism to create strings, either one must assume they exist ab initio, or one must conclude that string cosmologies require inflation. The non-perturbative potential dealt with here appears not to give the needed inflationary epoch.
1. Introduction. In $(2,2)$ string compactifications, where $(2,2)$ stands for $N = 2$ left-moving as well as $N = 2$ right-moving world-sheet supersymmetry, the vacuum is supersymmetric and contains a set of massless fields – moduli $T_i$ – which have no potential, i.e. $V(T_i) \equiv 0$, to all orders in string loops. Thus, perturbatively, there is a large degeneracy of string vacua, since any vacuum expectation value of moduli corresponds to the vacuum solution. On the other hand it is known that non-perturbative stringy effects, like gaugino condensation and axionic string instantons, give rise to the non-perturbative superpotential.

In the case of the modulus $T$ associated with the internal size of the compactified space for the so-called flat background compactifications (e.g., orbifolds, self-dual lattice constructions, fermionic constructions) the generalized target space duality is characterized by non-compact discrete group $PSL(2, \mathbb{Z}) = SL(2, \mathbb{Z})/\mathbb{Z}_2$ specified by

$$T \rightarrow \frac{aT - ib}{icT + d}, \quad ad - bc = 1, \quad \{a, b, c, d\} \in \mathbb{Z}.$$ (1)

Note, the dimensionless field $T$ is written as $T = R^2/\alpha' + iA$, where $R$ is the radius of compactification, $A$ corresponds to the internal axion field and $\alpha' (= 32\pi/g^2M_{pl}^2)$ is the string tension. Here, $g$ the gauge coupling (as defined in grand unified theories) and $M_{pl}$ is the Planck mass.

If one assumes that the generalized target space duality is preserved even non-perturbatively, the form of the non-perturbative superpotential is very restrictive. The fact that this is an exact symmetry of string theory even at the level of non-perturbative effects is supported by genus-one threshold calculations, which in turn specify the form of the gaugino condensate, and thus restrict the form of the non-perturbatively induced potential to a specific, restricted form of the duality invariant potential. In the case of the general duality invariant potentials, i.e., invariant under non-compact discrete symmetry $PSL(2, \mathbb{Z})$, this leads in general to stable supersymmetric domain walls. Such domain walls correspond to a new class of domain walls interpolating between non-degenerate minima of the matter potential with intriguing gravitational effects.

In this note we would like to address the cosmological consequences of a specific type of duality invariant potential, i.e., the one which is due to the gaugino condensation in the hidden sector of string theory. In the orbifold-type compactifications, due to the threshold corrections, such a duality invariant potential is of a very specific type.

Even when taking into account the Green-Schwartz mechanism to cancel anomalies of the underlying non-linear $\sigma$ model there is in general a duality invariant non-zero potential; e.g. it appears in the orbifold compactifications which have fixed two-tori in the twisted sector of the theory, like $Z_4$ orbifold. We will confine our study to such types of potentials.
The interaction of the modulus field $T$ is described in terms of $N = 1$ supergravity Lagrangian. The field lives in the fundamental domain $D$ determined by the modular group $\mathbb{C}/\text{PSL}(2,\mathbb{Z})$. (see fig. 1). We take the example of a duality invariant potential associated with one modulus field $T$ which corresponds to an overall size of the $T^6$ torus of compactified space. This is a general enough example which is to exhibit the qualitative features of gaugino induced non-perturbative potential. In this case the Kähler potential and the superpotential for the dilaton field $S$ and the modulus field $T$ can then be written in the form $^{[9,17]}$:

$$K = -\frac{1}{\kappa} \ln (S + S^*) - \frac{3}{\kappa} \ln (T + T^*)$$

$$W = \Omega(S) \eta(T)^{-6}$$

(2)

with $\kappa \equiv 8\pi G_N$. Here, $\eta(T)$ is the Dedekind eta function, a modular form of weight 1/2. Note, that $\eta$ is regular everywhere in the fundamental domain, and falls off as $e^{-1/12\pi T}$ as $T \to \infty$. The coefficient $\Omega(S)$ is in the case of the one hidden gauge group without matter of the type $M_C e^{24\pi^2/\ln S}$. $\Omega(S)$ sets the scale of the potential for the $T$ field and its exponential dependence on the $S$-dilaton field indicates that the superpotential is due to the non-perturbative physics. $b_0$ is related to the one-loop $N = 1$ (negative) beta function $\beta$ of the gauge group responsible for gaugino condensation as $\beta = b_0 g^3/16\pi^2$. Note, $ReS = 1/g^2$ where $g$ is the gauge coupling at the scale $M_C = gM_{pl}C$ where, constant $C$ depends on the subtraction scheme used $^{[20]}$. In $\overline{\text{DR}}$ scheme $C=0.043$. For $E_8$ hidden gauge group the scale $\Omega(S)^{1/3}$ is $\mathcal{O}(10^{16})$GeV.$^{\dagger}$

The scalar Lagrangian is of the form $^{\S}$:

$$e^{-1}L = -\frac{1}{2\kappa} R + K_{T\bar{T}} g^{\mu\nu} \partial_\mu T \partial_\nu \bar{T} - e^{\kappa K} (K^{T\bar{T}} |D_TW|^2 - 3\kappa |W|^2)$$

(3)

where $e = |\det g_{\mu\nu}|^{1/4}$, and $D_TW \equiv e^{-\kappa K} (\partial_T e^{\kappa K} W)$. For the Kähler potential and superpotential in (2) the Lagrangian(3)(with the fixed value of $S$) takes the form

$$e^{-1}L = -\frac{1}{2\kappa} R + \frac{3g^{\mu\nu} \partial_\mu T \partial_\nu T^*}{\kappa(T + T^*)^2}$$

$\dagger$ In the following we shall also neglect corrections due to one-loop corrections.$^{[16]}$

$\S$ In string theory the origin of a mechanism that fixes the dilaton field $S$ is not known. One possibility is the multiple gaugino condensation $^{[21,18]}$. We assume that $S$ is fixed to be $1/g^2$, with the gauge constant $g$ having perturbative value $< \mathcal{O}(1)$ which in turn determines scale associated with $\Omega(S)$. In further discussion we fix the value of $S$ and suppress the dynamics associated with the $S$ field, including the issue of supersymmetry breaking in this sector.

$\S$ We use $(+, -, -, -)$ space-time signature.
\[-\frac{\kappa|\Omega(S)|^2}{(T + T^*)^3(S + S^*)|\eta(T)|^{12}} \left[ \frac{(T + T^*)^2}{3} \left( \frac{3}{2\pi} \hat{G}_2(T, T^*) \right)^2 - 3 \right]. \tag{4}
\]

where \(\hat{G}_2 = -4\pi \partial_T \eta/\eta - 2\pi/(T + T^*)\) is the Eisenstein function of weight 2\textsuperscript{[19]}.

Note that without superpotential the moduli sector has a continuous non-compact symmetry \(SU(1, 1) \equiv SL(2, \mathbb{R})\). The non-perturbatively induced superpotential in (4) breaks the continuous symmetry down to its maximal discrete subgroup \(PSL(2, \mathbb{Z})\). The potential (see Fig. 3) in (4) has one minimum in the fundamental domain at \(T = 1.2\) which at the same time breaks supersymmetry. This turns out to be a generic property\textsuperscript{[9, 18]} for a gaugino induced non-perturbative potential in string vacua with \(PSL(2, \mathbb{Z})\) invariance.

2. Domain Walls Connected to Stringy Cosmic Strings. The underlying \(PSL(2, \mathbb{Z})\) symmetry of the theory implies\textsuperscript{[22, 8, 10]} that there could be domain wall solutions interpolating between such degenerate vacua of different fundamental domains.

The first class of domain walls is associated with the symmetry transformation \(T \rightarrow T + i\), \textit{i.e.}\ the discrete Peccei-Quinn symmetry. Thus the domain walls interpolates between minima with \(T = 1.2\) and \(T = 1.2 + i\). Note, that this transformation is a genuine symmetry of the underlying string vacua; \textit{e.g.}, the states of the underlying string vacua remain intact under such a symmetry transformation.

The nature of this type of domain wall is closely related to the domain wall that exists for the QCD induced potential of the Peccei-Quinn axion \(\theta\) when there is only one quark flavor. In this case \(V = \Lambda_{QCD}^4(1 + \cos \theta)\). In general, spontaneously broken global \(U(1)\) Peccei-Quinn symmetry is non-linearly realized through a pseudo-Goldstone boson, the invisible axion \(\theta = \{0, 2\pi\}\). Non-perturbative QCD effects through the axial anomaly break explicitly \(U(1)\) symmetry down to \(Z_{N_f}\), by generating an effective potential proportional to \(1 + \cos N_f \theta\). This potential leads to domain wall solutions\textsuperscript{[13]} with \(N_f\) walls meeting at the axionic strings\textsuperscript{[12]}.

In the case of \(N_f = 1\) there is still a domain wall, interpolating between \(\theta = 0\) and \(\theta = 2\pi\). The domain wall is bounded by an axionic string which emerged at the first stage of symmetry breaking of the global \(U(1)\) Peccei-Quinn symmetry.

Very analogous situation is taking place in our case; the role of the stringy axion field is played by the imaginary part of the \(T\) field; the domain wall interpolates between \(T = 1.2\) and \(T = 1.2 + i\); \textit{i.e.}\ inspite of only one minimum in the fundamental domain, there is still a domain wall just like in the case of axionic domain walls.

The analogy with axionic domain walls can be carried ever further. Such a
stringy domain wall is bound by a type of stringy cosmic strings.[23]

They can form at the first stage, when the *continuous* non-compact symmetry $SL(2, \mathbb{R})$ is spontaneously broken, *i.e.* at this stage there is only the kinetic energy term present while the non-perturbative potential due to gaugino condensation is not been turned on, yet (see eq.(4)).

In Ref.23 two types of such stringy cosmic strings were classified. They are obtained by mapping the $T$ field configuration on the $(x, y)$ spatial coordinates as:

$$j(T(x + iy)) = \sum_{n=0}^{n_{\text{max}}} a_n [(x + iy)/\sqrt{\kappa}]^n,$$

where $n$ are non-negative integers, $a_n$ are arbitrary constants (presumably of order one), and $j$ is modularly invariant function.[19] Note, that $j \to e^{2\pi T}$ as $T \to \infty$, thus at the spatial infinity ($x + iy \equiv r e^{i\phi}$ with $r \to \infty$) the $2\pi ImT$ is mapped into the spatial angle $2\pi n_{\text{max}}$ and $ReT$ blows up, *i.e.* decompactification takes place. Thus, as $T \to T + i$, the deficit angle is $\phi = 2\pi/n_{\text{max}}$. In addition, $j$ function has a triple zero at $T = e^{i\pi/6}$ and near $T = 1$ has an expansion $j = 1728 + a(T - 1)^2$.[19] This implies, that such a string has a core-like structure at two points, which are in the region $r = \mathcal{O}(\sqrt{\kappa})$, *i.e.* they correspond to $T = 1$ and $T = e^{i\pi/6}[23]$. The second type of domain walls corresponds to the map $j(T(x + iy)) = \sum_{n=0}^{n_{\text{max}}} a_n [(x + iy)/\sqrt{\kappa}]^n$. In this case the string core structure is at spatial infinity and decompactification takes place at $r = \mathcal{O}(\sqrt{\kappa})$. Note that the type of stringy domain walls associated with the symmetry transformation $T \to T + i$ can be bound only by the first type of stringy cosmic strings (see Fig. 3).

The second type of domain walls are associated with $T \to 1/T$, *i.e.* the generator of the non-compact symmetry transformation of $PSL(2, \mathbb{Z})$. This domain wall interpolates between the minimum at $T = 1.2$ and $T = 1/1.2$. It is analogous to the domain walls associated with $Z_2$ symmetry. It is at first puzzling that there would be such a domain wall, after all, the points in the $T$ plane are related by the $T \to 1/T$ symmetry. However, points associated with $T \to 1/T$ transformation can be probed since they correspond to a different theory (with heavy winding modes becoming light and vice versa) which happens to be equivalent to the original theory.

The second type ($T \to 1/T$) of domain walls are different in nature as opposed to the first type ($T \to T + i$) of domain walls. In particular, the second type of domain walls do not seem to be bound by stringy cosmic strings of the type described above.

The features of non-perturbatively induced potentials associated with the non-compact discrete symmetry for the radial moduli of other $(2,2)$ string vacua are

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[¶] In the case of supersymmetric stringy domain walls,[10], which appear in a class of general $PSL(2, \mathbb{Z})$ invariant potentials, the whole configuration of the domain wall bounded by the stringy cosmic string preserves supersymmetry.
generic. In general, the non-compact discrete symmetry group associated with
the radial modulus field has two types of symmetry generators: there is an analog of
the discrete Peccei-Quinn symmetry, i.e. \( T_k \rightarrow T_k + c \), with \( c \) being an integer or
half integer, as well as an analog of the non-compact symmetry (associated with the
symmetry of small and large radius of compactification). Thus, the two types of
domain walls discussed above seem to be a generic property of the gaugino induced
non-perturbative potential which respects the non-compact discrete symmetry of a
particular (2,2) string vacuum and possesses only one minimum in the fundamental
domain.

3. Cosmological Implications. In an expanding hot big bang cosmology the
walls form when the non-perturbative potential turns on. The scale of the domain
walls depends on the scale \( \Omega(S) \), i.e. the scale at which gaugino condensation
takes place. This scale could be as low as \( O(\text{TeV}) \) in the case of multiple gaugino
condensations \[21,18\]. In the case of hidden large gauge groups without matter a more
plausible scale is \( O(10^{16}\text{GeV}) \), which is favored in the case of hidden \( E_8 \) gauge
group. The domain walls produced by a phase transition such as this would have
degenerate vacua on either side and would quickly come to dominate the universe.
Thus, the walls have to go away in order to be consistent with observations.

In ordinary field theory such domain walls can disappear in two ways: one
possibility is inflation \[25\], another is chopping by a brownian network of cosmic
strings that appeared at an earlier phase transition \[26\]. It is the second possibility
which we will examine in the remainder of this section.

The difference between the standard picture and the present case is that the
stringy cosmic strings are inherently stringy; there is no phase transition in the
four dimensional effective field theory at which the strings appear. Rather, if
they were to exist they would have to have come directly from the spontaneous
compactification at the string epoch above \( M_C = 0.043 \ g \ M_{Pl} \), i.e. the scale below
which the physics of the effective four dimensional field theory takes over.

If so, then a problem for incorporating stringy cosmic strings in a cosmological
scenario is that there is apparently no natural choice for the string configuration
upon coming out of the Planck era. Recall that in cosmic string theories of galaxy
formation the strings are assumed, with some numerical and analytical support, to
evolve along a scaling solution in which the fraction of the energy density in strings
to that of the background is constant. On scales outside the horizon the network
executes a random brownian walk, its initial conditions set up at the time of a
phase transition according to the usual Kibble mechanism \[27\]. Note that the usual
Kibble mechanism only applies to strings which are created by a VEV appearing
at a phase transition. Just having fundamental strings in thermal equilibrium will

\* See, e.g., the study of properties of non-compact discrete symmetry for another string
vacuum \[24\].
not give a network of strings stretching beyond the horizon. In the case of stringy cosmic strings there is no such phase transition. We thus have no basis on which we can conclude that there would be a network of strings extending beyond the horizon. This problem is due to our lack of knowledge of superstring dynamics.

In spite of our ignorance of a mechanism by which they can be produced, let us assume that there is a Brownian string network present at the time that the non-perturbative potential turns on. The domain walls are then created bounded by strings. In our case such domain walls are bounded by stringy cosmic strings of the first type, only. E.g., \( j(T(x + iy)) = a_1(x + iy) + a_0 \). In this case cosmic string maps \( 2\pi ImT \to \phi \) as \( r \to \infty \). This is thus the type of a string that is attached to one end of the large domain wall. The energy per unit length of this string is\(^{[23]}\) in our case: \( \mu = 3/\kappa \times 2\pi/12 \). In addition, the core of such a string is more complicated as discussed above (see Fig. 3). Namely, the domain wall is bound by a string with the core like structure corresponding to two points: the \( Z_3 \) symmetric point at \( T = e^{i\pi/6} \) and \( Z_2 \) symmetric point \( T = 1 \). Also, the thickness of the domain wall \( (\mathcal{O}(\kappa |\Omega(S)|^2)^{-1/4}) \) can be much larger compared to the size of the string \( (\mathcal{O}(\sqrt{\kappa})) \).

Now suppose that walls form at a scale of \( \sigma \sim \mathcal{O}(10^{15}) \) GeV, where \( \sigma \) is the energy per unit area of the wall. We assume that a typical radius of curvature of a string at time \( t \) is less than \( t \), as is the case for standard cosmic strings. Then by the time \( t \sim (\mu/\sigma) \), where \( \mu \) is the string energy per unit length, the walls will dominate the dynamics of the system. However, because of the string network there will be no infinite domain walls. The walls-bounded-by-string quickly vanish by chopping into ever smaller pieces. Even if we make the extreme assumption that they fall into periodic motion, their lifetimes are limited by gravitational radiation to be less than \( \tau \sim 1/\sigma \), and the walls never come to dominate the universe\(^{[28]}\).

3. **Inflation?** If the absence of a standard Kibble mechanism to explain the existence of an initial string network means that the walls form unconnected to strings, then the only way to eliminate them is inflation. The non-perturbative potential in(4) for the \( T \) field could in principle allow\(^{[29]}\) for the inflationary epoch of the \( T \) field. Namely the potential shares common features with the pseudo-Goldstone potential \( \Lambda^4[\cos(\phi/f) + 1] \) with \( \Lambda = \mathcal{O}(10^{16}) \) GeV and \( f = \mathcal{O}(M_{pl}) \), which was shown\(^{[30]}\) to account for a version of inflation which is argued to be less finely-tuned than others. * In the case when \( f \geq M_{pl} \), inflation takes place for a wide range of intial field configurations, i.e. for \( \phi \to [0, \mathcal{O}(1)] \), while at the same time producing sufficient density fluctuations \( \delta \rho/\rho = \mathcal{O}(10^{-4}) \).

In our case, the major difference is that the field \( T \) is complex, and thus it rolls in the two-dimensional space. However, the potential is close to satisfying the

\* Within the study of instantons in a theory with the antisymmetric tensor, such constraint can also be met\(^{[31]}\).
two basic constraints. First, its overall scale is governed by the scale of gaugino condensation. From (4) one obtains \( \Lambda = [\kappa |\Omega(S)|^2]^{1/4} \), which in the case of hidden \( E_8 \) gauge group corresponds to the energy scale \( 10^{15} \text{GeV} \). In order to determine the scale \( f \), which governs the rate at which the potential changes, one has to normalize the kinetic energy in (4); this implies that the dimensionful field \( \phi \) is related to the dimensionless field \( T \) as \( \phi = T/\sqrt{\kappa} \). On the other hand the adjacent minima, are separated by a change in \( T \) of order one, e.g. along the imaginary direction \( T = 1.2 \) and \( T = 1.2 + i \). This in turn implies that the scale \( f = \mathcal{O}(1/(2\pi\sqrt{\kappa})) \).

In the fundamental domain \( \mathcal{D} \) a variation of the potential from its maxima and saddle points to its global minimum takes place for the region \( \text{Im}T = [0, 1/2] \) and \( \text{Re}T = [1, 1.2] \) (see Fig. 2). In our case, although one naturally expects \( f \sim M_{pl} \), when the field is properly rescaled we find that \( f \) is an order of magnitude too small for inflation to proceed via this truncation of the full theory. We must therefore have inflation arise from dynamics not considered in the present article. Another problem of the concrete example based on \( PSL(2, \mathbb{Z}) \) symmetry is the fact that the supergravity potential in (4) has a negative cosmological constant. However, the issue of ensuring zero cosmological constant is a question hardly addressed within any basic theory.

To conclude we would like to stress again that certain features of non-perturbatively induced potentials for the radial moduli of other (2,2) string vacua are generic. In general, the non-compact discrete symmetry group has two types of symmetry generators: there is an analog of the discrete Peccei-Quinn symmetry, i.e. \( T_k \rightarrow T_k + c \), with \( c \) being an integer or half integer, as well as an analog of the non-compact symmetry (associated with the symmetry of small and large radius of compactification). For such moduli the non-perturbative potential, which preserves the non-compact discrete symmetry and possesses only one minimum in the fundamental domain, in general has the properties discussed above. Without a mechanism to explain the existence of a cosmic string network we are led to the general conclusion such superstring universes must have inflation.

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REFERENCES

1. M. Dine and N. Seiberg, Nucl. Phys. B301, 357 (1988).

2. J. P. Derendinger, L. E. Ibáñez, and H. P. Nilles, Phys. Lett. 155B, 65 (1985); M. Dine, R. Rohm, N. Seiberg, and E. Witten, Phys. Lett. 156B, 55 (1985).

3. S.-J. Rey, *Axionic String Instantons and Their Low-Energy Implications*, Invited Talk at Tuscaloosa Workshop on Particle Theory and Superstrings, ed. L. Clavelli and B. Harm, World Scientific Pub., (November, 1989); Phys. Rev. D 43, 526 (1991).

4. S. Ferrara, D. Lüst, A. Shapere, and S. Theisen, Phys. Lett 225B, 363 (1989).

5. M. Cvetič, A. Font, L. E. Ibáñez, D. Lüst, and F. Quevedo, Nucl. Phys. B361, 194 (1991).

6. V. Kaplunovsky, Nucl. Phys. B307, 145 (1988).

7. L. Dixon, V. Kaplunovsky, and J. Louis, Nucl. Phys. B355, 649 (1991); J. Louis, *PASCOS 1991 Proceedings*, P. Nath ed., World Scientific 1991.

8. A. Font, L. E. Ibáñez, D. Lüst, and F. Quevedo, Phys. Lett. 245B, 401 (1990)

9. S. Ferrara, N. Magnoli, T. R. Taylor, and G. Veneziano, Phys. Lett. 245B, 409 (1990); P. Binetruy and M. K. Gaillard, Phys. Lett. 253B, 119 (1991); H. P. Nilles and M. Olechowski, Phys. Lett. 248B, 268 (1990).

10. M. Cvetič, S. J. Rey and F. Quevedo, Phys. Rev. Lett. 67, 1836 (1991).

11. M. Cvetič, S. Griffies and S.-J. Rey, *Static Domain Walls in N = 1 Supergravity*, UPR-474-T, YCTP-P43-91 (January 1992), Nucl. Phys. B, in press.

12. J. E. Kim, Phys. Rev. Lett. 43, 103 (1979); M. Dine, W. Fischler, and M. Srednicki, Phys. Lett. 104B, 199 (1981); and Nucl. Phys. B189, 575 (1981); M. B. Wise, H. Georgi, and S. L. Glashow, Phys. Rev. Lett. 47, 402 (1981); A good review is by J. E. Kim, Phys. Rep. 150, 1 (1987).

13. P. Sikivie, Phys. Rev. Lett. 48, 1156 (1982); G. Lazarides and Q. Shafi, Phys. Lett. 115B (1982) 21; For a review, see A. Vilenkin, Phys. Rep. 121, 263 (1985).

14. M. Cvetič and S. Griffies, *Gravitational Effects in Supersymmetric Domain Wall Backgrounds*, UPR-503-T (April 1992), Phys. Lett. B, in press.
15. G. Cardoso and B. Ovrut, Nucl. Phys. 369, 351 (1992), Proceedings of Strings and Symmetries, 1991, Stony Brook, World Scientific (P. Van Nieuwenhuizen et al. eds.) and Coordinate and Kähler σ-Model Anomalies and their Cancellation in the String Effective Field Theories, UPR-502-T (May 1992).

16. J.-P. Derendinger, S. Ferrara, C. Kounnas and F. Zwirner, CERN preprint CERN-TH-6004/91-REV and Phys. Lett. 271B, 307 (1991).

17. D. Lüst and C. Muños, Duality Invariant Gaugino Condensation and One-Loop Corrected Kähler Potentials in String Theory, CERN-TH 6358/91 (December 1991).

18. B. de Carlos, J. Casas and C. Muñoz, Supersymmetry Breaking and Determination of the Unification Gauge Coupling Constant in String Theory, CERN-TH-6346/92 (April 1992).

19. B. Schoeneberg, Elliptic Modular Functions, Springer, Berlin-Heidelberg (1970); J. Lehner, Discontinuous Groups and Automorphic Functions, ed. by the American Mathematical Society, (1964).

20. V. Kaplunovsky, Nucl. Phys. 307B, 145 (1988).

21. N. Krasnikov, Phys. Lett. 193B, 37 (1987); L. Dixon, V. Kaplunovsky, and M. Peskin, unpublished; L. Dixon, Proceedings of Rice Meeting of the APS DPF B. Bonner and H. Miettinen, eds. (World Scientific, Singapore, 1990); J. Casas, Z. Lalak, Muñoz, and G. Ross, Nucl. Phys. B347, 243 (1990).

22. A. Giveon, E. Rabinovici and G. Veneziano, Nucl. Phys. B322, 167 (1989).

23. B. Greene, A. Shapere, C. Vafa and S.-T. Yau, Nucl. Phys. B337, 1 (1990).

24. P. Candelas, X. De la Ossa, P. Green and, L. Parkes, Phys. Lett. 258B, 118 (1991).

25. A. H. Guth, Phys.Rev. D23, 347 (1981); A. D. Linde, Phys. Lett. 108B, 389 (1982); A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. 48, 1220 (1982).

26. T. W. B. Kibble, G. Lazarides and Q. Shafi, Phys. Rev. D26, 435 (1982); A. Vilenkin and A. E. Everett, Phys. Rev. Lett. 48, 1867 (1982); A. E. Everett and A. Vilenkin, Nucl. Phys. B207, 43 (1982).

27. T. W. B. Kibble, J. Physics A9, 3320 (1976); Physics Reports 67, 183 (1980).

28. A. Vilenkin, Phys. Rep. 121, 263 (1985).

29. M. Cvetič, F. Quevedo and S.-J. Rey, unpublished.

30. K. Freese, J. Frieman and A. Olinto, Phys. Rev. Lett. 65, 3233 (1990).

31. B. Ovrut and S. Thomas, Phys. Lett. 267B, 227 (1991) and Instantons in Antisymmetric Tensor Theories in Four-Dimensions, UPR-0465-T (1991).
Figure Captions

Fig.1 Fundamental domain of the T field. The dot \( T = e^{i\pi/6} \) denotes the \( Z_3 \) symmetric point and the cross \( T = 1 \) denotes the \( Z_2 \) symmetric point.

FIG.2 The scalar potential for the T field in units of \( \kappa|\Omega(S)|^2 \). The absolute minimum lies on the real axis at \( T = 1.2 \). The figure is taken from \([8,5]\).

FIG.3 Stringy cosmic string connected to a domain wall. The locations of the stringy core at \( r = 0 \) \( (T = e^{i\pi/6}) \) and at \( re^{i\phi} = 1728\sqrt{\kappa} \) \( (T = 1) \) corresponds to the stringy cosmic string with the map \( j(T) = re^{i\phi}/\sqrt{\kappa} \).