

CERN-TH.7255/94

THE ROLE OF
THE SUPERSTRING DILATON
IN COSMOLOGY AND PARTICLE PHYSICS

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ABSTRACT

Superstring theory predicts the existence of a scalar field, the dilaton. I review some basic features of the dilaton interactions and explain their possible consequences in cosmology and particle physics.

^{*)} Talk given at XXIXth Rencontres de Moriond, Electroweak Interactions and Unified Theories, March 12-19 1994, Meribel, France.

CERN-TH.7255/94
May 1994

1. Introduction

In this talk I concentrate on the role of the superstring dilaton in determining cosmological evolution in the early universe and some particle physics properties at energies well below the Planck scale. The idea is that even without solving the full problem of non-perturbative dynamics in string theory some general conclusions can still be drawn. The unique role of the dilaton in this context has been recognized previously by many authors (See e.g. [1] and references therein).

I review some simple and universal features of the superstring dilaton and show that, when supplemented with information about the general nature of the dilaton potential, they have rather well defined consequences for the cosmological evolution in the early universe. For particle physics more detailed models of supersymmetry breaking are necessary. Some additional fields beside the dilaton have to be included. So far, there does not exist a satisfactory class of models.

2. The Superstring Dilaton

In this section I review some properties of the superstring dilaton. Most of them are known for quite some time[1, 2]. The superstring dilaton is a scalar field that appears in all variants of (super)string theory. Its universal existence and some characteristics of its interactions can be traced directly to basic symmetries of the theory.

One way of understanding the dilaton appearance is as a consequence of 2-dimensional invariance under (super)conformal symmetry on the world-sheet. The propagation of (super)strings in space-time backgrounds can be described in terms of a two-dimensional generalized σ -model. (Super)Conformal invariance of this theory is a condition that determines in which backgrounds (super)strings can propagate consistently. The dilaton field, ϕ appears as a specific coupling in the Lagrangian density[3],

$$\mathcal{L}_{dil}^{(2)} = \sqrt{g}R^{(2)}\phi(X^\mu) \tag{1}$$

of the σ -model. Here, X^μ is the sigma-model fields and g , $R^{(2)}$ are the determinant of the 2-dimensional metric, and 2-dimensional curvature scalar, respectively. It turns out that even if this term is not included classically,

the condition of (super)conformal invariance always requires its generation at higher orders. Note that the integral $\int d^2z \sqrt{g} R^{(2)}$ is a topological number, proportional to the Euler character of the 2-d manifold.

Another way of understanding the dilaton appearance is as a consequence of 10-dimensional supersymmetry in space-time. If one starts with 10-d superstrings one finds that in order for the space-time supersymmetry transformations to close, the supergravity supermultiplet of the theory has to contain in addition to the metric (vielbein), an antisymmetric tensor field and a massless scalar field, the dilaton.

Yet another way to view the dilaton is as a consequence of a 10-dimensional classical dilatation symmetry of the effective low energy field theory. The bosonic part of the Lagrangian density can be written as

$$\mathcal{L} = \frac{1}{16\pi\alpha'} \sqrt{-g} e^{-\phi} \left\{ R + \nabla\phi \cdot \nabla\phi - \frac{1}{3}H^2 - \frac{1}{4}\text{Tr}F^2 + \dots \right\} \quad (2)$$

where α' is the string tension and g is the gauge coupling parameter, F is the gauge field strength and H is the field strength of the antisymmetric tensor. The dots stand for other terms not specified here. Once a vacuum expectation value (VEV) is chosen the classical dilatation symmetry is spontaneously broken. In this context the dilaton can be thought of as the Goldstone boson of dilatation symmetry and hence the name dilaton. A large degeneracy exists in the absence of a mechanism that fixes dynamically the VEV of the dilaton.

The more interesting string theories are 4-dimensional string theories. The dilaton in 4-d string theory can be thought of as a part of the 10-d dilaton. The superstring dilaton in 4-d is best described in terms of a chiral superfield, S , of $N = 1$ supergravity. The bosonic part of $S = e^{-\phi} + ia$ contains as its real part the dilaton and addition to the dilaton another scalar field, the axion. The axion is a part of the 10-d antisymmetric tensor. The Lagrangian is then given in terms of the Kahler potential $K = -\ln(S + S^*)$, and the superpotential W . The axion field interactions are invariant under a shift symmetry inherited from the antisymmetric tensor field gauge symmetry. This symmetry is expected to be broken along with supersymmetry in low energy effective field theory.

Some basic and universal features of the dilaton interactions can be directly determined from the above characterization of the dilaton. From

Eq.(1) one sees that the zero-mode of the dilaton couples to the Euler character of the 2-d world-sheet manifold. The Euler character determines the genus of the 2-d manifold which in turns determines the number of loops in the corresponding string diagram. The quantity e^ϕ therefore serves as a string loop counting parameter. For the same reason the VEV of the dilaton is related to the string coupling parameter and gauge coupling parameter. This can be seen from Eq.(2) as well as the fact that the dilaton couples to all the fields in the theory with gravitational strength.

Additional important information about the dilaton interactions is obtained from models of supersymmetry (and shift symmetry) breaking . Typically[4], supersymmetry breaking is accompanied by the generation of a potential for the dilaton. One important requirement is that the potential vanishes at large negative values of ϕ , corresponding to the weak coupling region[2]. The dilaton potential is conveniently described as perturbative or non-perturbative. Perturbative potentials have exponential dependence on ϕ , $V_{pert.} \sim e^{(n-1)\gamma\phi}$ where n is the order of perturbation theory in which they are generated and γ is a fixed numerical parameter. Non-perturbative potentials have even stronger dependence on ϕ , $V_{non-pert.} \sim e^{-\alpha e^{-\phi}}$. The general shape of the potential and the way it approaches zero at weak coupling are thus fixed. The details of the shape of these potentials, such as the existence of minima or maxima etc. are model dependent. The most popular assumption is that some non-perturbative effects are responsible for the generation of the dilaton potential. The weakness of these effects may explain the hierarchy between the electroweak scale and the Planck scale.

3. The Role of the Superstring Dilaton in Cosmology

The cosmological evolution of the early universe is particularly affected when the VEV of the dilaton is not constant in space and time. The situations in which the VEV of the dilaton varies in space-time are either if dilaton is away from the minimum of its potential or if the dilaton potential vanishes altogether, or if temperatures and energies are so high that the potential becomes unimportant. In this “dilaton-roll” epoch, coupling parameters that are usually considered constants, such as Newton’s constant or the charge of the electron , vary with time or even in space. There are strict experimental limits on the amount that various coupling constants are

allowed vary in space-time today[5]. We therefore know that the dilaton-roll epoch has ended, most probably before the nucleosynthesis epoch[2, 6, 7]. The bounds on variation of coupling parameters in the very early universe are much more relaxed.

During the dilaton-roll epoch the presence of the dilaton changes the gravitation evolution equations from the regular Einstein equations to equations of the Brans-Dicke type. The questions one would like to answer about this epoch are: Is there a phase of inflationary evolution and if so how is it affected by the dilaton? How does the dilaton eventually settles down to the minimum of the potential. For generic dilaton potentials it is possible to obtain definite answers by solving the classical equations of motion, either analytically or numerically.

When the VEV of the dilaton is constant in space-time, the gravitational evolution equation are the same as Einstein's equations. This situation is expected to occur at later times when the dilaton field is sitting at the minimum of its potential¹. The value of the full dilaton potential at the minimum is the late time cosmological constant of the theory. A natural requirement is that this value vanishes to the appropriate accuracy. Another cosmological parameter associated with the evolution of the dilaton field in the vicinity of the minimum of its potential is the rate of dilaton particles production as the field oscillates coherently and settles down to the minimum of its potential[9]. The dilaton couples with gravitational strength and is generically light and may therefore induce changes in late time cosmological evolution or even cause deviations from the equivalence principle today[10]. These issues can be fully resolved only in a framework of a more detailed model of supersymmetry breaking.

The description of cosmological evolution in the dilaton-roll epoch is through a solution of the dilaton-gravity equations of motion. A class of interesting solutions are of the Friedmann-Robertson-Walker (FRW) type, in which the metric is isotropic, with vanishing spatial curvature and the dilaton depends only on time.

$$\begin{aligned} ds^2 &= -dt^2 + a^2(t)dx_i dx^i & i = 1, 2, 3 \\ \phi &= \phi(t) \end{aligned} \tag{3}$$

The Hubble parameter, H , is related to the scale factor, a in the usual way,

¹An alternative mechanism has been suggested[8]

$H \equiv \frac{\dot{a}}{a}$. It is possible to make a field redefinition from the string (Brans-Dicke) frame, as in Eq.(2), to the Einstein frame in which the curvature and dilaton are canonically normalized. The dilaton-gravity equations of motion in the Einstein frame are

$$\begin{aligned} \ddot{\phi} + 3H\dot{\phi} &= -\frac{dV(\phi)}{d\phi} \\ H^2 &= 8\pi\left(\frac{1}{2}\dot{\phi}^2 + V\right)/(3M_{Pl}^2) \end{aligned} \quad (4)$$

Without the potential this system of equations corresponds to a Brans-Dicke field with $\omega = -1$. Inflationary evolution requires a rather long phase of super-luminal expansion followed by a transition to a regular expanding FRW universe. The nature of the evolution is determined by the shape of the dilaton potential and initial conditions. In particular the important feature is the steepness of the potential. A useful example is the case of an exponential potential, $V(\phi) = V_0 e^{\beta\phi}$. In this case, the solution can be obtained analytically $\phi(t) = \phi_0 - \frac{2}{\beta} \ln \frac{t}{t_0}$ and $a(t) = a_0 \left(\frac{t}{t_0}\right)^{4\pi/\beta^2}$ for $t > t_0$. If the parameter β , that determines the steepness of the potential, is above a critical value $\beta_{crit} = 2\sqrt{\pi}$ no super-luminal expansion occurs.

The steepness of the dilaton potential is determined by the 2-d coupling in Eq.(1). For tree level potentials, the steepness exactly corresponds to β_{crit} . It allows an expansion at exactly the speed of light but not more [11, 6]. As explained, potentials that are induced at higher orders in perturbation theory are steeper than the tree level potential and therefore the scale factor expansion is suppressed compared to the tree level. Non-perturbative potentials are generically steeper than any of the perturbatively induced potentials and therefore in this case the expansion is even slower[7].

The dilaton couples to all fields in a multiplicative form. The rate of scale factor expansion is determined by competition on the conversion of potential energy into kinetic energy in which the field whose potential is the steepest wins. This means that potential-driven inflation is hard to come by and is generically absent in string theory in the dilaton-roll epoch.

An alternative cosmological evolution where an interplay between the dilaton and scale factor kinetic energies is important has been suggested[12, 13], examined[14] and described in more detail in another talk[15]. I will describe it briefly and state the conclusions. The analysis is based on the equations of motion of the same dilaton-gravity system in the original string

frame. These can be written as a system of two first order equations.

$$\begin{aligned}\dot{H} &= \pm H\sqrt{3H^2 + V} - \frac{1}{2}V' \\ \dot{\phi} &= 3H \pm \sqrt{3H^2 + V}\end{aligned}\tag{5}$$

The (\pm) signifies that either a $(+)$ or $(-)$ is chosen for both equations simultaneously. The solutions to the equations (5) belong to two branches, according to which sign is chosen. Another twofold ambiguity is related to the sign of H . The $(-)$ branch is related to the solutions that were considered previously. This branch can be joined smoothly to an ordinary FRW radiation-dominated expanding universe after the dilaton-roll epoch has ended. This is the reason why most analyses so far concentrated on this branch.

The $(+)$ branch has instead some unusual and quite remarkable properties. In the absence of any potential the solution for $\{H, \phi\}$ is now given by $H^{(+)} = \pm \frac{1}{\sqrt{3}} \frac{1}{t-t_0}$ and $\phi^{(+)} = \phi_0 + (\pm\sqrt{3} - 1) \ln(t_0 - t)$, $t < t_0$. This solution describes either accelerated expansion and evolution from a cold, flat and weakly coupled universe towards a hot, curved and strongly coupled one [13] or accelerated contraction and evolution towards weak coupling. In general, the effects of a potential on this branch are quite mild. The dilaton zooms through potential minima. It is impossible, in this branch, for the dilaton to sit at a minimum of its potential. Inflation, in this solution, is driven, not hampered, by the dilaton's kinetic term, thanks to the negative value (-1) of the BD ω parameter.

The possibility of branch changes in the weak curvature and weak coupling region was checked[14] with negative results. The apparent conclusion is that a transition from an accelerated inflation era to a standard FRW cosmology is not possible in the weak curvature regime for any sort of (semi) realistic dilaton potential. The remaining possibilities lie therefore in the strong curvature (large derivatives) regime, in which the full extent of string corrections should be felt. Final conclusions have to await finding a good model for this type of cosmological evolution, but at the moment it seems as a promising direction[16].

4. The Role of the SuperString Dilaton in Particle Physics

Particle physics properties of the low energy effective field theory are in general associated with the minimum of the dilaton potential and its vicinity. The VEV of the dilaton determines the gauge coupling parameter at a high scale. In addition the dilaton is a part of the “hidden sector” of supergravity (See e.g. [17]). It is a field whose only interactions with the observable fields, charged under the $SU(3) \times SU(2) \times U(1)$ gauge group, are of gravitational strength. As a rule more detailed information about the dilaton potential and the type of supersymmetry (and shift symmetry) breaking mechanism are necessary. Obviously, just specifying an exponential or a non-perturbative dependence of the potential is not enough. Both type of potentials do not have a minimum and this is inconsistent with the fact that dilaton has to have a constant VEV today.

A class of models that are described² in terms of a non-perturbative superpotential $W(S) = \sum a_i e^{-\alpha_i S}$ and a Kahler potential $K = -\ln(S + S^*)$ was proposed [18] developed[19] and analysed in detail[20]. In these models a_i and α_i are real numbers. The parameters α_i are all rather large, $\alpha_i \sim 15$. The sum is a finite small sum. I will use properties of this class of models to explain what type of additional information is required. The main characteristic feature of this class of superpotentials is that they are very steep, i.e. have large derivatives in the mild coupling region. A property that was already important in the discussion of cosmological evolution.

There are some basic conditions that the dilaton potential has to satisfy to be considered as a viable model for the hidden sector of Supergravity in the present context. The value of the potential at the minimum is part of the cosmological constant. It is only part because it does not include for example the contribution from the Higgs potential after electroweak symmetry breaking. The potential has to be chosen such that this value is set to zero to the correct accuracy. This does not solve the cosmological constant problem but allows one, in practice, to ignore the problem[21]. The expectation value of the dilaton at the minimum is related to the value of the gauge coupling parameter at some high scale. This value is expected to be small. The reasons for this prejudice are twofold. First, if one follows the running of the various gauge coupling parameters they seem to meet at a value which is still small

²All the formulae here are tree level results.

(See e.g. [22]). The other is that if the coupling is strong, one expects that the classical spectrum and interactions, which are the reason why string theory is interesting to begin with, will be changed appreciably. In addition, the simplifying assumption that the coupling parameter remained small through out the evolution is usually made.

It is tempting to assume that some universal conclusions may be drawn by studying the dilaton on its own[23]. This is not possible within the class of models previously described. To see this, let us assume for the moment that it is possible to consider just the dilaton and that the effects of all other fields can be represented through values of parameters in the dilaton potential and kinetic terms. The potential is given by

$$V(S) = (S + S^*)|F_S|^2 - \frac{3}{S + S^*}|W(S)|^2 \quad (6)$$

here S is the chiral superfield mentioned in Sect. 2 and $F_S = \partial_S W(S) - \frac{1}{S+S^*}W(S)$. The VEV of F_S determines whether supersymmetry is broken or not. Assume now the following properties for $W(S)$

- a) $|W(S)| \rightarrow 0$ for $S \rightarrow \infty$
- b) $|(S + S^*)\partial_S^{(n+1)}W/\partial_S^n W| \gg 1 \quad n = 0, 1, \dots$

in the weak coupling region, $S + S^* > 0$. Condition *b*) states that W has large derivatives in the weak coupling region. The condition for an extremum of the potential is $\frac{\partial V}{\partial S} = 0$,

$$(S + S^*)\partial_S^2 F_S^* - \frac{2}{S + S^*}F_S W^* = 0 \quad (7)$$

This equation can be satisfied in two different ways

i) $F_S = 0$

or

ii) $(S + S^*)^2|\partial_S^2 W| = 2|W|$ and $F_S \neq 0$

Because of condition *b*) above, it is only possible to satisfy Eq.(7) in case *i*) if $\partial_S W \sim 0$ and in case *ii*) if $\partial_S^2 W \sim 0$. Consider an extremum of type

ii). Condition *b*) can be used to analyse the matrix of second derivatives at the extremum. The conclusion is that the extremum is a saddle points or a maximum but not a minimum.

For a superpotential $W(S)$ that is a real function it is possible to analyse the situation completely also for type *i*) extrema[7], if they occur at a real value of S . These extrema are generically minima[20]. If $W \neq 0$ at the minimum, the cosmological constant is negative. Note that the fact that $F_S = 0$ for this type of minima and the fact that the cosmological constant is negative are correlated. If $W = 0$ at the minimum, S_{min} , as well, the cosmological constant vanishes. However, since $W(\infty) = 0$ and $W(S_{min}) = 0$ there is a a point in between where $\partial_S W = 0$. From Eq.(6) one sees that at this point $V < 0$ and therefore the global minimum of the potential in the weak coupling region is negative. The implications of this analysis for the gaugino condensation models was explained in detail elsewhere[7].

The conclusion is that one has to consider explicitly more fields in the hidden sector, or a different type of superpotentials or both. This means that particle physics properties will depend mainly on the structure of the additional fields and the nature of changes to the superpotential and not just on universal properties of the dilaton. This point of view was presented by other authors as well[24, 25]. A suggestion about the nature of additional fields in the spirit of no scale models has been recently put forward[26].

Acknowledgements

This talk is based in part on joint work with Paul Steinhardt and with Gabriele Veneziano. I would like to thank Costas Kounnas and Fabio Zwirner for discussions.

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