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DILATON COUPLINGS AT LARGE DISTANCES

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Abstract

The large-distance behaviour of dilaton couplings is studied within the framework of four-dimensional heterotic superstring theories. String loop corrections violate universality of dilaton and graviton couplings, increasing dilaton couplings at large distances. This implies a serious problem for string theories, unless a non-zero dilaton mass is generated.

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The presence of dilatons – massless, totally gauge-neutral scalars – seems to be inevitable in string theories. One generally expects that dilatons interact with matter no stronger than gravity, i.e. with couplings of order $\mathcal{O}(\kappa \equiv \sqrt{4\pi G_N})$; the number of dilatons, as well as detailed forms of their couplings, depend on particular models under consideration.

The physical existence of dilatons is very much constrained by the results of experiments designed to study deviations from Einstein's theory of general relativity. Typically, such experiments provide various bounds on the dilaton couplings, depending on the dilaton mass.[1]

The determination of the dilaton mass presents a very difficult problem, related to ultra-violet string physics, supersymmetry breaking etc. On the other hand, as argued in this letter, some general aspects of dilaton interactions, like their large-distance behaviour, do *not* depend on the ultra-violet physics and can be studied by using an effective field theory approximation. To be specific, we consider here only the minimal case of one dilaton field, within the framework of so-called four-dimensional heterotic superstring theories [2], however our analysis could be readily extended to more complicated dilaton systems.

In four-dimensional heterotic superstring theories, the dilaton field ϕ (more precisely, its exponential) belongs to the chiral supermultiplet S, whose interactions with matter and gauge fields are completely determined by the Kähler potential $J = -\log(S+\bar{S})+\dots$ [3]. Let (A,ψ) and (B,ξ) denote two chiral matter superfields of mass $m \ll 1/\kappa$, in the representations r and $\bar{r}$ of some generic gauge group, respectively.[2] The interaction Lagrangian for this dilaton - gauge - matter system is given by [3]:

$$\begin{aligned}\mathcal{L}_{\rm int} \;=\; & -\kappa\phi[m\psi\xi + m\bar{\psi}\bar{\xi} + 2m^2A^\dagger A + 2m^2BB^\dagger] \\ & -i\sqrt{2}g\kappa\phi[A^\dagger T^{(a)}\psi\lambda^{(a)} - BT^{(a)}\bar{\xi}\bar{\lambda}^{(a)} - \bar{\lambda}^{(a)}\bar{\psi}T^{(a)}A + \lambda^{(a)}\xi T^{(a)}B^\dagger] \\ & -g^2\kappa\phi[A^\dagger T^{(a)}A - BT^{(a)}B^\dagger]^2 \;+\; \frac{1}{2}\kappa\phi F^{(a)}_{\mu\nu}F^{(a)\mu\nu} \;+\; \mathcal{O}(\kappa^2)\,, \end{aligned} \tag{1}$$

where g denotes the gauge coupling constant and $T^{(a)}$ the gauge group generators in the representation r. In eq.(1) we used the notation of ref.[5].

The dilaton interactions of eq.(1) exhibit the property that one would generally expect

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[1]See e.g. ref.[1].

[2]Such small masses can be generated in some phenomenologically interesting superstring models by a string analogue of the Higgs mechanism [4].

for any massless neutral scalar: at the classical level, the zero-momentum coupling to a particle on its mass-shell is proportional to its rest mass. Since the dilaton under consideration originated from string theory, it is not surprising that its coupling constant is equal to the graviton coupling constant κ. Such a "strongly" coupled massless dilaton would lead to unacceptable deviations from Einstein's theory of general relativity; for instance, non-relativistic measurements of the gravitational constant would yield a result twice as large as the observation of light deflection from the sun. In the case of such a strong coupling, laboratory, geophysical and astronomical data exclude a dilaton mass lower than 10^{-4} eV [6]. Could radiative corrections change this conclusion?

Instead of using the fully-fledged string apparatus in the calculation of radiative corrections to the dilaton couplings, we shall first follow the effective field theory approach of ref.[7], to estimate the size of string-loop effects; to be specific, we discuss the one-loop effects, however our arguments apply to an arbitrary number of loops as well.

In order to determine the one-loop dilaton couplings to the order $\mathcal{O}(\kappa)$, we evaluate the kinetic energy corrections and the one-particle irreducible dilaton vertices. The kinetic energy corrections should be absorbed into rescalings of fields and masses, so that the one-loop effective Lagrangian, expressed in terms of these rescaled fields, acquires the canonical form of the kinetic energy terms. Finally, the dilaton vertices should be expressed in terms of these rescaled fields and masses.

Since dimension d operators in the Lagrangian are weighted by the factor κ^{d-4}, which is of the order of the respective power of the ultra-violet string momentum cut-off $1/\alpha' = g^2/2\kappa^2$, only the loops of light (i.e. of masses $m \ll 1/\kappa$) string excitations may give rise to large corrections, by yielding logarithms which become singular in the infra-red limit of vanishing masses or momenta. Moreover, a simple power-counting argument shows that the loops involving gravitational interactions are free of infra-red singularities to the order $\mathcal{O}(\kappa)$.[3] The only terms that require special attention are the matter mass terms with "unnaturally" small $m \ll 1/\kappa$, however in the supersymmetric theory under consideration their smallness is protected to all orders by chiral symmetries that are restored in the limit of $m \to 0$; some logarithmic corrections are expected, though. We thus conclude that only the loops involving light particles and gauge interactions may give rise to non-trivial radiative corrections, provided that some infra-red singularities are present. This justifies the use of quantum field theory, instead of a fully-fledged string apparatus, in the calculation of these large contributions.

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[3]Soft graviton singularities [8] are absent to this order.

As far as the technical aspect of this calculation is concerned, the simplest way to extract the leading logarithms from the Feynman diagrams is to use dimensional regularisation of ultra-violet divergences, by continuing momentum integrals to $D = 4 - 2\epsilon$ dimensions.[4] We subtract the ultra-violet ϵ^{-1} poles, and take the limit $\epsilon \to 0$. At this point, the scale μ of the leading logarithmic terms, of order $g^2 \log(\frac{m}{\mu})^2/16\pi^2 \sim \mathcal{O}(1)$, should be identified with the string ultra-violet cut-off: $\mu \sim 1/\kappa$. One important comment is in order here. Radiative corrections induce some interactions which are absent at the tree-level, for instance $\phi\partial_\mu A^\dagger \partial^\mu A$ etc. These terms, as well as corrections to the terms already present in eq.(1), are individually gauge-dependent, however as expected, this gauge-dependence cancels out in on-mass-shell amplitudes, i.e. after using equations of motion.

As already mentioned in the introduction, we are mainly interested in the large distance behaviour of the dilaton couplings, therefore we restrict our attention to dilaton couplings at zero momentum transfer. We obtain the following form of the effective on-mass-shell scattering Lagrangian, quadratic in the matter fields:

$$\mathcal{L}^{\text{eff}}_{\text{quadr}} = -K\kappa\phi[m_R\psi\xi + m_R\bar{\psi}\bar{\xi} + 2m_R^2 A^\dagger A + 2m_R^2 BB^\dagger] \ + \ \mathcal{O}(\kappa^2)\,, \tag{2}$$

where now all matter fields correspond to rescaled canonical fields, with canonical kinetic terms and inertial mass:

$$m_R^{1-\text{loop}} = m[1 - \frac{g^2 C(r)}{8\pi^2} \log(\kappa^2 m^2)]\,. \tag{3}$$

In eq.(3), $C(r)$ denotes the quadratic Casimir operator for the representation r. The factor K of eq.(2), which parametrises deviations of the dilaton coupling from the graviton coupling constant, is given by:

$$K^{1-\text{loop}} = 1 - \frac{g^2 C(r)}{4\pi^2} \log(\kappa^2 m^2)\,. \tag{4}$$

In order to extend our result, eqs.(2)-(4), beyond the one-loop level, and to perform summation of the leading logarithmic terms, it is convenient to make use of the sigma-model approach to string theory. Let $\tilde{G}_{\mu\nu}$ and $\tilde{\phi}$ denote the background metric and the background dilaton field, respectively, of the non-linear string sigma-model. A generic string h-loop contribution to the effective Lagrangian is of the form [9]:

$$\alpha' \Delta\mathcal{L}^{\text{eff}} = \sqrt{-\tilde{G}}\, \tilde{O} e^{2(h-1)\tilde{\phi}} g^{2h-2} Z_O^{(h)}\,, \tag{5}$$

[4] No divergences other than logarithmic are encountered in the two- and three-point Green's functions under consideration; these are exactly the infra-red singular logarithms we want to extract.

where $\tilde{O}$ is an operator without the constant $\tilde{\phi}$ mode. In four dimensions, the sigma-model metric and dilaton are related to the physical metric $G_{\mu\nu}$ and dilaton ϕ (of the S-matrix approach) [9] by the following equations:

$$\tilde{G}_{\mu\nu} = G_{\mu\nu} e^{2\kappa\phi}, \qquad \tilde{\phi} = \kappa\phi. \tag{6}$$

For an operator $\tilde{O}$ of conformal weight w, we obtain the following coupling of the zero-momentum physical dilaton:

$$\alpha' \Delta\mathcal{L}^{\rm eff} = \sqrt{-G}\, O e^{2(1+w)\kappa\phi} g^{2h-2} Z_O^{(h)} e^{2h\kappa\phi}. \tag{7}$$

Note that the h-loop dilaton coupling to the operator O is completely determined by the h-loop contribution $Z_O^{(h)}$ to the operator O.

Since the four-dimensional string theory under consideration is supersymmetric, the higher-loop corrections to the factor K, eq.(2), can be determined by studying the dilaton coupling to the scalar field A. The effective Lagrangian for this dilaton - scalar system is given by:

$$\begin{aligned}
\alpha'\mathcal{L}^{\rm eff}(A,\phi) &= -\partial_\mu A^\dagger \partial^\mu A[\sum_{h=0}^{\infty} g^{2h-2} Z_k^{(h)} e^{2h\kappa\phi}] - m^2 A^\dagger A e^{2\kappa\phi}[\sum_{h=0}^{\infty} g^{2h-2} Z_m^{(h)} e^{2h\kappa\phi}] \\
&= -\partial_\mu A^\dagger \partial^\mu A[\sum_{h=0}^{\infty} g^{2h-2} Z_k^{(h)}] - m^2 A^\dagger A[\sum_{h=0}^{\infty} g^{2h-2} Z_m^{(h)}] \\
&\quad - 2\kappa\phi\{\, \partial_\mu A^\dagger \partial^\mu A[\sum_{h=0}^{\infty} h g^{2h-2} Z_k^{(h)}] + m^2 A^\dagger A[\sum_{h=0}^{\infty} (h+1) g^{2h-2} Z_m^{(h)}]\,\} \\
&\quad + \mathcal{O}(\kappa^2).
\end{aligned} \tag{8}$$

We now perform the rescaling of the scalar field A, to bring its kinetic energy to the canonical form. In terms of this rescaled field, the Lagrangian of eq.(8) is given by:

$$\begin{aligned}
\mathcal{L}^{\rm eff}(A,\phi) &= -\partial_\mu A^\dagger \partial^\mu A - m_R^2 A^\dagger A \\
&\quad - 2\kappa\phi\{\, \partial_\mu A^\dagger \partial^\mu A\,(g^2 \frac{\partial}{\partial g^2} \log Z_k) + m_R^2 A^\dagger A\,(1 + g^2 \frac{\partial}{\partial g^2} \log Z_m)\,\},
\end{aligned} \tag{9}$$

where the inertial mass is given by:

$$m_R^2 = m^2 \frac{Z_m}{Z_k}, \tag{10}$$

$$Z_k = \sum_{h=0}^{\infty} g^{2h} Z_k^{(h)}, \qquad Z_m = \sum_{h=0}^{\infty} g^{2h} Z_m^{(h)}. \tag{11}$$

After using the equations of motion for the scalar field A, we obtain the following result for the factor K, defined in eq.(2):

$$K \;=\; 1+g^2\frac{\partial}{\partial g^2}\log(\frac{Z_m}{Z_k}) \;=\; 1+g^2\frac{\partial}{\partial g^2}\log m_R^2\,. \tag{12}$$

The leading logarithmic terms in the mass rescaling factor $Z_m Z_k^{-1}$, eq.(10), can be summed up to all orders in the gauge coupling constant by using the standard renormalisation group techniques. From the one-loop result of eq.(3), we obtain:

$$\frac{Z_m}{Z_k} = \left[\frac{g^2(m^2)}{g^2}\right]^{4C(r)/\beta_0}, \tag{13}$$

where the running coupling constant is defined by the usual leading logarithmic relation:

$$\frac{1}{g^2(m^2)} = \frac{1}{g^2} + \frac{\beta_0}{16\pi^2}\log(\kappa^2 m^2)\,. \tag{14}$$

By substituting eq.(13) into eq.(12), we obtain the following result for the factor K in the leading logarithmic approximation:

$$K^{\text{llog}} = 1 - \frac{g^2(m_R^2)C(r)}{4\pi^2}\log(\kappa^2 m_R^2)\,. \tag{15}$$

Thus in the leading logarithmic approximation, the inclusion of higher-loop effects amounts to the replacement of the coupling constant g^2 in the one-loop result of eq.(4) by the running coupling constant corresponding to the energy scale set by the rest mass of particles interacting with the dilaton. Note that the dilaton coupling is always *larger* than the graviton coupling constant, for any gauge group or representation r.

There is another way of deriving eq.(12), which sheds some light on the dilaton interactions in confining theories, for instance on the dilaton - nucleon interactions. Eq.(5) shows that to all orders in the string loop expansion parameter g (gauge coupling constant):

$$\mathcal{L}^{\text{eff}}(\tilde\phi,\tilde G_{\mu\nu}) \;=\; \mathcal{L}^{\text{eff}}(\tilde\phi+\log g,\tilde G_{\mu\nu}) \;=\; \mathcal{L}^{\text{eff}}(\kappa\phi, G_{\mu\nu}e^{2\kappa\phi})\,, \tag{16}$$

so that:

$$\frac{\partial}{\partial\tilde\phi}\mathcal{L}^{\text{eff}} \;=\; 2g^2\frac{\partial}{\partial g^2}\mathcal{L}^{\text{eff}}\,. \tag{17}$$

The coupling of the zero-momentum physical dilaton ϕ is given by:

$$\frac{\partial}{\partial\phi}\mathcal{L}^{\text{eff}}(\phi=0) \;=\; 2\kappa(-G^{\mu\nu}\frac{\partial}{\partial G^{\mu\nu}} + g^2\frac{\partial}{\partial g^2})\mathcal{L}^{\text{eff}}(\phi=0)\,. \tag{18}$$

For a canonically normalised field A, the most general quadratic form consistent with general covariance, is given by:

$$\mathcal{L}_{\rm quadr}(\phi=0,G_{\mu\nu},A) \;=\; -\sqrt{-G}(G^{\mu\nu}\partial_\mu A^\dagger\partial_\nu A + m_R^2 A^\dagger A)\,. \tag{19}$$

After applying eq.(18) and using the equations of motion, we obtain:

$$\mathcal{L}^{\rm eff}_{\rm quadr}(\phi,A) \;=\; -2\kappa(m_R^2 + g^2\frac{\partial m_R^2}{\partial g^2})\phi A^\dagger A \;=\; -2\kappa(1+g^2\frac{\partial}{\partial g^2}\log m_R^2)m_R^2\phi A^\dagger A\,, \tag{20}$$

in agreement with the result of eq.(12) for the factor K. The same result can be derived explicitly for fermions, without using supersymmetry. Of course, in a supersymmetric theory the factor K is equal for fermions and bosons of equal masses.

It is clear from our derivation that eq.(20) follows from: i) general covariance and ii) particular dependence of the effective action on the background dilaton field: $\mathcal{L}^{\rm eff}(\tilde{\phi}) = \mathcal{L}^{\rm eff}(\tilde{\phi}+\log g)$. It is generally expected that non-perturbative effects do not spoil general covariance. Also, it seems reasonable to assume that low-energy phenomena do not affect the Planck-scale relation between the coupling constant and the dilaton background. In other words, we expect that, even in a confining theory like QCD, the dilaton couples to the rest mass m_R of composite particles, e.g. nucleons, with the factor K given in eq.(12). In a confining theory, the typical particle mass is of order of the strong interaction scale $\Lambda \sim \kappa^{-1}\exp(-8\pi^2/\beta_0 g^2)$ [7], therefore we expect:

$$K \;\approx\; 1+\frac{16\pi^2}{\beta_0 g^2} \;\gg\; 1\,. \tag{21}$$

Similar considerations can be repeated for the graviton couplings, with the conclusions that: i) the graviton coupling κ remains unaffected by radiative corrections, ii) the inertial mass m_R is equal to the gravitational mass, i.e. the mass that couples to spin-two gravitons. The equivalence principle is *not* violated in the graviton sector, as expected from general covariance.

We conclude that the universality of dilaton and graviton couplings does not hold beyond the tree-level; as seen from eqs.(15) and (21), at low energies the dilatons couple stronger than the gravitons. The effective non-relativistic coupling, induced by the superposition of one-dilaton and one-graviton exchanges is equal to $(1+K)\kappa$. In the presence of dilaton interactions, the equivalence principle *is* violated: non-relativistic matter interacts with couplings different from massless gauge bosons that couple to gravitons only, with coupling equal to κ. Moreover, since the dilaton coupling contains the factor K, eq.(12), which depends on the rest mass, even the weak equivalence principle is violated.

In summary, we studied radiative effects in dilaton interactions with the conclusions that at large distances dilatons couple stronger than gravitons, and violate the equivalence principle. We confirm that without addressing the dilaton mass generation problem, it is not possible to avoid serious phenomenological problems in string theories with massless dilatons.

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