Properties of the dilaton

T. E. Clark, C. N. Leung, and S. T. Love
Department of Physics, Purdue University, West Lafayette, Indiana 47907
(Received 16 July 1986)

Gauge models which possess a hierarchy of mass scales may exhibit a pseudo-Goldstone boson, the dilaton, of the scale symmetry. Properties of the dilaton are studied in a technicolor model, where the technifermion chiral condensation scale differs from the confinement scale. The dilaton exhibits properties very similar to the elementary Higgs boson of the standard model. It is distinguished from the Higgs boson by its Goldstone nature, which is manifested in an enhanced multidilaton production. Estimates for the dilaton mass and lifetime are given and various production mechanisms are discussed. Prospects for detecting the dilaton in its associated production with W or Z and in W and Z decays are stressed.

I. INTRODUCTION

In a recent Letter,[1] some of us have suggested that the spontaneous breaking of chiral symmetry in certain gauge models may also trigger the spontaneous breaking of an approximate scale symmetry, resulting in the appearance of a pseudo-Goldstone boson, the dilaton. This would be the situation if the chiral-symmetry breaking occurs when the explicit breaking of the scale symmetry due to quantum fluctuations is not a dominant effect. Such may be the case when fermions belonging to higher-dimensional representations of the gauge group are present in the theory. Indeed, results from numerical studies in lattice gauge theory indicate that the scale of chiral condensation for these fermions is relatively short compared to the confinement scale.[2] This is anticipated since the chiral condensation scale is characterized by the requirement that the effective fermion coupling $C_2(f)g^2(\mu)$ reaches a certain critical value. Here $g(\mu)$ is the running coupling constant of the underlying gauge theory and $C_2(f)$ is the quadratic Casimir invariant of the fermion representation. For a sufficiently large $C_2(f)$, spontaneous chiral-symmetry breaking can occur in the asymptotically free regime where $g(\mu)$ varies only logarithmically with energy. The resulting explicit breaking of the scale symmetry is thus a small effect compared to the large spontaneous breaking of scale symmetry associated with the chiral condensation. In such cases, the theory should be reliably approximated by an effective theory with a fixed but critical coupling constant. This fixed coupling theory now possesses an exact scale invariance in the chiral-symmetry limit. When the chiral symmetry is spontaneously broken, so is the scale symmetry. As a result, the Goldstone boson of scale symmetry, the dilaton, appears as a massless fermion-antifermion scalar bound state. Actually the coupling constant is not fixed and the scale symmetry is explicitly broken by the scale anomaly which reflects the running of the coupling constant. Therefore the dilaton should appear as a pseudo-Goldstone boson with a mass of order the scale at which the explicit scale symmetry breaking is important, which is roughly the confinement scale of the gauge theory.

In this paper we study the phenomenological consequences resulting from the assumed existence of the dilaton.[3] To be definite, we study the effects of the dilaton in the context of a specific model. However, many qualitative features of our results are actually more general and depend only on the fact that the dilaton is a Nambu-Goldstone boson of dilatation symmetry.[4]

II. DESCRIPTION OF THE MODEL

The model we consider is a QCD-like technicolor (TC) model[5] with a weak $SU(2)_L$ doublet of technifermions transforming as some irreducible representation of the TC gauge group which is presumably larger than the fundamental representation. It is assumed that the chiral condensation of the technifermions forms at a scale μ_{TC} which is much larger than the TC confinement scale Λ_{TC}. The confinement scale is defined to be the energy scale at which the gauge coupling constant equals one, viz., $\alpha_{TC}(\Lambda_{TC})=1$, while the scale μ_{TC} is characterized by the requirement that $C_2(f)\alpha_{TC}(\mu_{TC})=\alpha_{\text{critical}}\simeq 1$, where $C_2(f)$ is the quadratic Casimir invariant of the technifermion representation. The technifermions transform under $SU(3)_C\times SU(2)_L\times U(1)_Y$ as

$$\begin{aligned} T_L &= \begin{bmatrix} X_L \\ Y_L \end{bmatrix} \in (1,2,0) , \\ X_R &\in (1,1,\tfrac{1}{2}) , \\ Y_R &\in (1,1,-\tfrac{1}{2}) . \end{aligned} \tag{1}$$

The weak-hypercharge assignment is chosen such that the model is anomaly-free. This leads to electric charges of $\frac{1}{2}$ and $-\frac{1}{2}$ for X and Y.

The electroweak gauge symmetry is spontaneously broken to electromagnetism at the scale $\mu_{TC}\gg\Lambda_{TC}$ by the dynamical condensate

$$|\langle \bar{X}X+\bar{Y}Y\rangle| \sim {\mu_{TC}}^3 . \tag{2}$$

A dynamical Higgs mechanism ensues and the three Goldstone bosons of the $SU(2)_L\times U(1)_Y\to U(1)_{EM}$ symmetry breakdown comprise the dominant contributions to the longitudinal components of the massive $W^{\pm}$ and Z^0 vector bosons. The mass of the W boson is obtained as

$$M_W = \frac{eF}{2\sin\theta_w} , \tag{3}$$

while the phenomenologically successful relation $M_W = M_Z \cos\theta_w$ emerges automatically. Here θ_w is the usual weak mixing angle and $F \sim \mu_{\rm TC}$ is the Goldstone-boson decay constant defined by

$$\langle 0 \mid J^i_\mu \mid S^j(k) \rangle = iF\delta_{ij}k_\mu , \tag{4}$$

where J^i_μ ($i=1,2,3$) are the spontaneously broken SU(2) currents and S^i are the associated Goldstone bosons. This decay constant effectively replaces the vacuum expectation value of the elementary scalar field in the standard-model Higgs mechanism. Its value can be determined from the measured values of M_W, $\sin^2\theta_w$, and α yielding $F = 246$ GeV.

In addition, the chiral condensate (2) also breaks the scale symmetry spontaneously. Since, by assumption $\mu_{\rm TC} \gg \Lambda_{\rm TC}$, the TC gauge coupling constant varies only logarithmically with energy near the scale $\mu_{\rm TC}$. This results in only a weak explicit breaking of the scale symmetry as opposed to the large spontaneous breaking resulting from the chiral condensate (2). Consequently, a fourth Goldstone boson, the dilaton (D), will appear in the spectrum as a scalar technifermion-antitechnifermion bound state. Since the explicit breaking of the scale symmetry becomes important when the coupling constant becomes strong (i.e., when $\mu \sim \Lambda_{\rm TC}$), the dilaton will not be an exactly massless Goldstone boson, but rather, a pseudo-Goldstone boson with a mass of order $\Lambda_{\rm TC}$. The coupling of the dilaton to matter is characterized by a decay constant F_D defined by

$$\langle 0 \mid J^D_\mu \mid D(k) \rangle = iF_D k_\mu , \tag{5}$$

where J^D_μ is the dilatation current. In general, one expects F_D to be roughly the same magnitude as F.

The interactions of the dilaton with the usual fermions (quarks and leptons) and the gauge bosons can be described by an effective Lagrangian which nonlinearly realizes the scale symmetry. The form of this effective Lagrangian is presented in the Appendix. The dilaton behaves very much like the elementary Higgs boson. (This is a general consequence of the dilaton's coupling to the trace of the energy-momentum tensor and does not depend on the specific model considered here.[6]) For example, it couples to the usual quarks and leptons as

$$\mathcal{L}_{Df\bar{f}} = \sum_f \frac{m_f}{F_D} D\bar{f}f , \tag{6}$$

where m_f denotes the mass of the fermion f. Since $F_D \sim F = 246$ GeV, the strength of these couplings are comparable to that of the Higgs boson. In general, one distinction between technicolor models and models with elementary scalar fields is that there is no Higgs-type scalar (which couples via the mass) in the bound-state spectrum of the technifermions. However, when the technifermion condensation occurs at a scale at which the effect of the explicit scale symmetry breaking is suppressed, a Higgs-type dilaton may appear. The dilaton distinguishes itself from the Higgs boson in its Goldstone nature, which is apparent from its couplings to the longitudinal W and Z bosons and manifested in an unusually large cross section for multidilaton production at high energies, as we will see later. Furthermore, contrary to the Higgs boson whose mass can be anywhere between roughly 7 GeV (Ref. 7) and about 1 TeV (Ref. 8), the dilaton is necessarily light compared to the scale of electroweak symmetry breaking.

A novel feature of the present class of TC models is the possibility of automatic suppression[9] of flavor-changing neutral currents arising from the extended technicolor interactions.[10] In the usual TC theories, it is necessary to postulate a new interaction, namely, extended technicolor (ETC), in order to generate fermion masses. For an asymptotically free TC theory, a typical fermion mass is given by $m_f \sim \mu_{\rm TC}{}^3/M_{\rm ETC}{}^2$ (up to logarithms),[11] where $M_{\rm ETC}$ is the mass of the ETC gauge bosons assumed to be heavy ($M_{\rm ETC} \gg \mu_{\rm TC}$). However, the ETC interactions also generate flavor-changing neutral couplings via the effective four-fermion interactions $M_{\rm ETC}{}^{-2}\bar{f}f\bar{f}f$. Present experimental limits on flavor-changing neutral currents require $M_{\rm ETC} \geq 10^3$ TeV. This would give too small a fermion mass. In TC theories where the explicit scale symmetry breaking is suppressed when the technifermions condense, it has been speculated that the ultraviolet behavior be controlled by a nontrivial fixed point.[3,9,12] Consequently, the technifermion self-energy has the large-momentum behavior $\Sigma(p) \sim \mu_{\rm TC}{}^2/p$. This leads to a fermion mass given by $m_f \sim \mu_{\rm TC}{}^2/M_{\rm ETC}$. For the same m_f and $\mu_{\rm TC}$, $M_{\rm ETC}$ is a factor $\sqrt{\mu_{\rm TC}/m_f}$ larger than that in the asymptotically free TC models. Thus, a realistic fermion mass can be generated by the ETC interactions without violating the constraints imposed by flavor-changing neutral couplings.

III. MASS AND LIFETIME ESTIMATES

As stated earlier the dilaton develops its mass when the explicit breaking of the scale symmetry becomes important. This occurs when the running coupling constant $\alpha_{\rm TC}$ becomes of order one. The mass of the dilaton m_D is thus expected to be of the order of the TC confinement scale $\Lambda_{\rm TC}$. Unfortunately a reliable estimate of m_D is not available.[12] However, since $\Lambda_{\rm TC}$ is assumed to be small compared to $\mu_{\rm TC} \sim 246$ GeV, we expect $m_D \ll 246$ GeV and probably no larger than about 100 GeV; such a value is not unreasonable in light of the results of Ref. 2. On the other hand, $\Lambda_{\rm TC}$ should be at least as large as $\Lambda_{\rm QCD}$ and probably larger. We therefore take m_D to be in the range[13]

$$\text{few hundred MeV} \lesssim m_D \lesssim 50 \text{ GeV} . \tag{7}$$

There exists experimental limits on the mass of a Higgs-type scalar particle. None, however, excludes the mass range of Eq. (7). Studies in nuclear transitions and neutron-nucleus scattering experiments[14] yield the bound $m_D > 18$ MeV. Absence of the decays $K^\pm \to \pi^\pm l^+ l^-$ ($l = e$ or μ) can be used to set the limit $m_D \gtrsim m_K - m_\pi \simeq 350$ MeV (Ref. 15). The validity of this limit has recently been challenged.[16] (See, however, Ref. 17.) Searches in heavy-quarkonia decays (J/ψ or $\Upsilon \to \gamma + D$) suffer from large QCD corrections[18] and do not yield useful limits at present (see the discussion in Sec. IV).

The lifetime of the dilaton depends on two parameters, its mass m_D and the decay constant F_D. As does the Higgs boson, the dilaton tends to decay into the heaviest possible fermion pair. The width for the dilaton to decay into a fermion pair is

$$\Gamma(D\to f\bar{f})=N\frac{m_D}{8\pi}\left[\frac{m_f}{F_D}\right]^2\left[1-\frac{4m_f^2}{m_D^2}\right]^{3/2}$$
$$\equiv\tilde{\Gamma}\left[\frac{m_D}{50\text{ GeV}}\right]\left[\frac{250\text{ GeV}}{F_D}\right]^2, \tag{8}$$

where N is the color factor (3 for quarks, 1 for leptons); and the lifetime of the dilaton is

$$\tau_D=\left[\sum_{\substack{f\\ m_D>2m_f}}\Gamma(D\to f\bar{f})\right]^{-1}$$
$$\equiv\tilde{\tau}\left[\frac{50\text{ GeV}}{m_D}\right]\left[\frac{F_D}{250\text{ GeV}}\right]^2. \tag{9}$$

Note that τ_D is inversely proportional to m_D and directly proportional to F_D^2. Depending on the exact value of F_D, the lifetime of the dilaton can be either shorter or longer than that of a Higgs boson with the same mass.

We summarize in Table I the lifetime and dominant decay modes of the dilaton for various values of its mass. For a light dilaton ($m_D\lesssim 1$ GeV), it decays predominantly to $\mu^+\mu^-$ and $\pi\pi$. We anticipate that the $\pi\pi$ decay rate should be at least as large as the $\mu^+\mu^-$ decay rate and thus the naive estimate

$$\frac{\Gamma(D\to\pi\pi)}{\Gamma(D\to\mu^+\mu^-)}\simeq\frac{3(m_u^2+m_d^2)}{m_\mu^2}, \tag{10}$$

where m_u and m_d are current-quark masses, appears to be an underestimate.[19] Since a reliable estimate for the two-pion decay rate is lacking, we provide only an order-of-magnitude estimate in Table I. For a heavier dilaton, more decay channels open up. Its lifetime varies from about 10^{-17} sec for $m_D=1$ GeV to about 10^{-22} sec for $m_D=50$ GeV, assuming only three generations with $m_t>25$ GeV and $F_D=250$ GeV. Sample branching ratios for $m_D=10$ GeV and $m_D=50$ GeV are shown in Table II.

The dilaton can also decay into gauge bosons. Decays to W^+W^- and Z^0Z^0 are not energetically accessible. The decay to $\gamma\gamma$ is only important if $m_D<2m_e$. For the mass range of interest, the $\gamma\gamma$ decay rate is negligible. For m_D greater than about 4 GeV (which is assumed to be above the glueball threshold), the decay $D\to gg$ is allowed. This decay proceeds via the diagram in Fig. 1 and is dominated by heavy quarks in the fermion loop. The rate is given by

$$\Gamma(D\to gg)=\frac{1}{8\pi}\left[\frac{\alpha_s(m_D)}{3\pi}n_h\right]^2\frac{m_D^3}{F_D^2}, \tag{11}$$

where n_h is the number of quarks heavier than the dilaton and α_s is evaluated at m_D. Hence,

$$\frac{\Gamma(D\to gg)}{\Gamma(D\to\tau^+\tau^-)}=\left[n_h\frac{\alpha_s(m_D)}{3\pi}\right]^2\left[\frac{m_D}{m_\tau}\right]^2$$
$$\times\left[1-\frac{4m_\tau^2}{m_D^2}\right]^{-3/2} \tag{12}$$

which is typically a few percent. The decay $D\to gg$ is thus relatively unimportant having a branching ratio of order 1% or less. It will, however, become more important if there is a fourth generation of quarks.

IV. PRODUCTION AND DETECTION

Since the dilaton couples to quarks and leptons just like the elementary Higgs boson (with F_D replacing the vacuum expectation value), its production mechanism will be very similar to that of the Higgs boson. As F_D is expected to be about 250 GeV, the production cross sections will also be comparable to those of the Higgs boson with one important exception. This is that the dilaton has an unusually large multiparticle production rate at high energies due to its Goldstone nature. Ways for producing and detecting a Higgs boson with mass in the range (7) have been studied in Ref. 6 and reviewed in Ref. 20. We discuss here those mechanisms which are more promising for the dilaton.

TABLE I. Partial widths and lifetime of the dilaton [$\tilde{\Gamma}$ and $\tilde{\tau}$ are defined in Eqs. (8) and (9)]. The t quark is assumed to be heavier than 25 GeV/c^2.

m_D (GeV/c^2)	Dominant decay modes	$\tilde{\Gamma}$	$\tilde{\tau}$
0.3–1.0	$\mu^+\mu^-$	0.36 keV	
	(e^+e^-)	(8.5×10^{-3} eV)	10^{-19}–10^{-18} sec
	$\pi^+\pi^-,\pi^0\pi^0$	$\sim$1 keV	
1–4	$s\bar{s}$	2.2 keV	$\sim2\times10^{-19}$ sec
	$\hookrightarrow K^+K^-,K^0\bar{K}^0,\ldots$		
4–10	$\tau^+\tau^-$	100 keV	2×10^{-21} sec
	$c\bar{c}$	215 keV	
10–50	$b\bar{b}$	2.2 MeV	2.6×10^{-22} sec

TABLE II. Sample branching ratios of the dilaton.

Decay modes	Branching ratios (%) $m_D = 10$ GeV	$m_D = 50$ GeV
$\mu^+\mu^-$	0.1	0.01
$\pi\pi$	0.1–0.2	0.01–0.02
$s\bar{s}$	0.7	0.09
$\tau^+\tau^-$	31	4
$c\bar{c}$	68	8.5
$b\bar{b}$		87

A. Production from the decays of other particles

1. Heavy-quarkonia decays

The fact that the dilaton couples more strongly to heavier quarks makes it attractive to search for the dilaton in heavy-quarkonia decays: $V \to D + \gamma$ ($V = c\bar{c}$, $b\bar{b}$, or $t\bar{t}$). A light dilaton ($m_D \simeq 1$ GeV) should be searched for in the decays of ψ or Υ, while a heavier dilaton can be seen in t-quarkonium decays. Similar to the Higgs boson, the rate of the radiative decay to the dilaton is given by

$$r \equiv \frac{\Gamma(V \to D + \gamma)}{\Gamma(V \to e^+e^-)} = \frac{1}{8\pi\alpha}\left[\frac{m_V}{F_D}\right]^2\left[1 - \frac{m_D^2}{m_V^2}\right], \tag{13}$$

where m_V is the mass of the quarkonium state V. Unfortunately, there is a large $O(\alpha_s)$ QCD correction which tends to lower the rate.[18] For instance, the ratio r for the Υ is estimated to be reduced by a factor of about 2 (Ref. 21). It is not known whether higher-order QCD corrections would further reduce the rate. Consequently, present knowledge of ψ and Υ decays do not provide useful limits on m_D or F_D and future monochromatic photon searches need improvement in sensitivity. If t quarkonium is discovered, it will be an ideal laboratory to search for the dilaton. Not only is the rate (13) larger due to the large $m_{t\bar{t}}$, but also the QCD radiative correction is smaller because α_s is smaller at $q^2 = m_{t\bar{t}}^2$.

2. Gauge-boson decays

The properties of the W and Z bosons will be extensively studied both in e^+e^- machines and hadron colliders in the next few years. In particular, the Z boson will be studied at resonance using the Stanford Linear Collider. It is therefore of utmost interest to search for the dilaton in the decays of W and Z. Indeed the prospects of finding the dilaton in these searches is quite favorable if the dilaton is sufficiently light ($m_D \lesssim 20$ GeV).

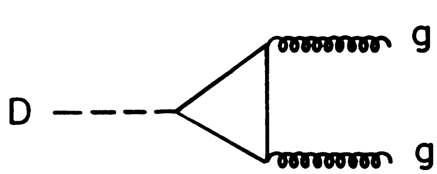

FIG. 1. Decay of the dilaton into two gluons.

a. $Z \to D + f\bar{f}$ *or* $W \to D + f\bar{f}'$. These decays proceed via the diagrams in Fig. 2 with the dominant contribution coming from Fig. 2(a). The coupling of the dilaton to the W and Z bosons can be gleaned from $\mathcal{L}_{\rm GD}$ in (A14) as

$$\mathcal{L}_{DWW} = \frac{2M_W^2}{F_D} D W_\mu^+ W^{-\mu}, \tag{14}$$

$$\mathcal{L}_{DZZ} = \frac{M_Z^2}{F_D} D Z_\mu Z^\mu. \tag{15}$$

These are similar to the couplings of a Higgs boson to the W and Z with F_D replacing $(G_F\sqrt{2})^{-1/2}$. There are also contributions from $\mathcal{L}_A$ in (A9) to $\mathcal{L}_{DWW}$ and $\mathcal{L}_{DZZ}$. These are however, suppressed by $\beta(g)/g$.

The rate for $Z(W) \to D + f\bar{f}(\bar{f}')$ can be directly obtained from the analogous process involving a Higgs boson with the same mass using the relation

$$\Gamma[Z(W) \to D + f\bar{f}(\bar{f}')] = \frac{1}{G_F\sqrt{2}F_D^2}\Gamma[Z(W) \to H + f\bar{f}(\bar{f}')]. \tag{16}$$

We thus find the branching ratio $B(Z \to D + \mu^+\mu^-) \simeq 3 \times 10^{-5}(250\ \text{GeV}/F_D)^2$ for $m_D = 10$ GeV. This reduces to about $10^{-6}(250\ \text{GeV}/F_D)^2$ for $m_D = 50$ GeV. With 10^6 or more Z events, a dilaton with mass 20 GeV or less can definitely be seen. The signature is in fact rather clean. If the dilaton does not decay in the detector, one would see missing mass recoiling against the muon pair. If the dilaton decays inside the detector, then according to Table II, a considerable fraction of the decays would be to $\tau^+\tau^-$. One would then see a τ pair opposite to a μ pair (or four muons in the case of a lighter dilaton).

The decay $W \to D + l\nu_l$ ($l = e$, μ, or τ) has a larger branching ratio, $B(W \to D + l\nu_l) \simeq \text{few} \times 10^{-4}(250\ \text{GeV}/F_D)^2$, but is harder to detect because of the missing neutrino. Nevertheless, there is a good chance one can

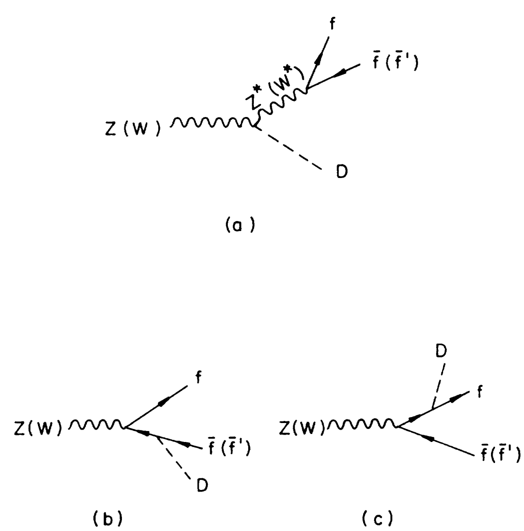

FIG. 2. Decays of Z (or W) to $D + f\bar{f}$ (or $f\bar{f}'$).

find the dilaton (particularly a light one) in W decays. A light dilaton which decays to a muon pair would contribute to like-sign dimuon events in W decays.

b. $Z \to D + \gamma$. There is a direct dilaton-Z-photon coupling in the effective Lagrangian which is obtained from (A9) as

$$\mathscr{L}_{DZ\gamma} = -\frac{D}{4F_D}\sin 2\theta_w \left[\frac{\beta_2^{\mathrm{TF}}(g_2)}{g_2} - \frac{\beta_1^{\mathrm{TF}}(g_1)}{g_1}\right] Z_{\mu\nu}F^{\mu\nu}, \tag{17}$$

where $Z_{\mu\nu} = \partial_\mu Z_\nu - \partial_\nu Z_\mu$ and $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$. The β functions depend on the technicolor representation of the technifermions. For moderate dimension representations this contribution to the $Z \to D + \gamma$ decay rate is much smaller than the contribution from the W-boson loop depicted in Fig. 3(a). [There is also a contribution from fermion loops in Fig. 3(b), but it is negligible for $m_f \ll M_W$.] The W-boson loop contribution can be estimated from the corresponding $Z \to H + \gamma$ decay rate[22] as

$$\frac{\Gamma_W(Z \to D + \gamma)}{\Gamma(Z \to \mu^+\mu^-)} \simeq 8 \times 10^{-5}\left[1 - \frac{m_D{}^2}{M_Z{}^2}\right]^3 \times \left[1 + 0.16\frac{m_D{}^2}{M_Z{}^2}\right]\left[\frac{250\ \mathrm{GeV}}{F_D}\right]^2 . \tag{18}$$

Using the branching ratio $B(Z \to \mu^+\mu^-) \simeq 3\%$, we then obtain, for $m_D \ll M_Z$,

$$B(Z \to D + \gamma) \simeq 3 \times 10^{-6}\left[\frac{250\ \mathrm{GeV}}{F_D}\right]^2 . \tag{19}$$

It is thus more favorable to search for light dilatons in $Z \to D + f\bar{f}$ decays. On the other hand, for heavier dilatons, the rate for $Z \to D + f\bar{f}$ decreases rapidly and becomes comparable to that for $Z \to D + \gamma$. The decay $Z \to D + \gamma$ would then be more attractive because of its clean signature.

B. Production in association with W and Z bosons

Another way of producing the dilaton is by associated production with W or Z in either e^+e^- collisions ($e^+e^- \to Z^* \to Z + D$) or $p\bar{p}$ collisions [$p\bar{p} \to Z^*(W^*) \to Z(W) + D$]. This is of imminent interest because the dilaton is relatively light and thus will be accessible to experiments in the immediate future. The production rate is governed by the couplings (14) and (15) and hence related to the corresponding associated production of a Higgs boson of the same mass by (σ denotes cross sections)

$$\sigma_D = \frac{1}{G_F\sqrt{2}F_D{}^2}\sigma_H . \tag{20}$$

If F_D is less than 246 GeV, it will be easier to produce a dilaton than to produce a Higgs boson. Detailed cross sections for Higgs production can be found in the review by Ansel'm, Ural'tsev, and Khoze.[20] We will not reproduce them here.

C. Production by gluon fusion

Because of its small couplings to light quarks, it is easier to produce the dilaton in $\bar{p}p$ (or pp) collisions by gluon fusion, analogous to the gluon-fusion mechanism for Higgs production.[23] The cross section is again related to that for a Higgs boson with the same mass by (20). Detailed cross sections for Higgs production can be found in Refs. 20 and 23.

D. Multidilaton production

As stated earlier the dilaton has a relatively large multiparticle production rate at high energies. This is a result of the Goldstone nature of the dilaton and is analogous to the enhanced multiple production of longitudinal W and Z at high energies.[24] Consider the two-dilaton production as shown in Fig. 4. The dilaton self-coupling is given by $\mathscr{L}_D$ in (A10). Because of the derivative coupling, the q^{-2} in the dilaton propagator is canceled by the q^2 in the three-dilaton vertex (we are considering situations where $q^2 \gg m_D{}^2$). The phase-space integral produces a factor proportional to s. Hence, aside from numerical factors of π, the ratio of the two-dilaton production rate to the one-dilaton production rate is

$$\frac{\sigma_{2D}}{\sigma_{1D}} \sim \frac{s}{F_D{}^2} . \tag{21}$$

For $\sqrt{s} \gtrsim F_D$, two dilatons can be produced almost as often as one dilaton. For a light dilaton, this leads to rather distinctive signatures: four or more charged leptons (μ or τ).

V. CONCLUSION

We have presented a case study of the dilaton in the context of a technicolor model. An effective Lagrangian for the dilaton interactions is given. The dilaton behaves in many ways like the elementary Higgs boson of the stan-

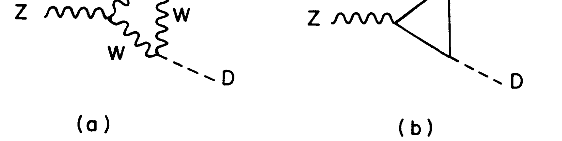

FIG. 3. One-loop contributions to the decay $Z \to D + \gamma$.

FIG. 4. Multidilaton production.

dard model. They differ in that the dilaton is a pseudo-Goldstone boson and is expected to be relatively light. The Goldstone character also implies an enhanced multidilaton production at high energies analogous to the enhanced multiple production of longitudinal W and Z in strongly coupled Higgs models. The mass and lifetime of the dilaton are estimated and various ways to detect it are discussed. A lighter dilaton should be searched for in the decays of W and Z and heavy quarkonia, while a heavier one can be produced in association with W or Z.

The dilaton can in principle mix with a 0^{++} glueball or a 0^{++} techniglueball. We have not studied the effects of such mixings. One should be able to distinguish the dilaton from a glueball by their leptonic branching ratios. A 2-GeV dilaton will decay to $\mu^+\mu^-$ about 10% of the time whereas the branching fraction of a glueball to $\mu^+\mu^-$ is $O(\alpha^2)$. In our simple model, the techniglueball is rather stable and can be distinguished from the dilaton by its long lifetime.

ACKNOWLEDGMENTS

We thank W. A. Bardeen for useful correspondence and J. Pantaleone and J. L. Rosner for stimulating discussions. S.T.L. thanks the SLAC Theory Group for their hospitality during the completion of the manuscript and the DOE for financial support under the Outstanding Junior Investigator Program. This work was supported in part by the U.S. Department of Energy.

APPENDIX

We derive in this appendix the effective Lagrangian which governs the physics below the technifermion chiral condensate scale, μ_{TC}, but above the technicolor confinement scale Λ_{TC}, which is assumed to be small compared to μ_{TC}. The technifermions develop constituent masses of order μ_{TC} via the mechanism of dimensional transmutation and no longer act as relevant degrees of freedom below μ_{TC}. The technigluons, however, remain relevant physical degrees of freedom. Hence, the effective Lagrangian is constructed with the dilaton D, the Goldstone bosons S^i $(i=1,2,3)$ associated with the spontaneous breakdown of the technifermion chiral SU(2)$\times$SU(2) symmetry to a vectorial SU(2)$_V$ symmetry, the technigluon fields T^μ, the quarks and leptons of the standard model as well as the color and electroweak gauge bosons.

Rather than working in the unitary gauge in which the S^i disappear from the spectrum through a Higgs effect, we shall work in the R_ξ class of gauges. This will be useful for calculating processes that involve the longitudinal weak gauge bosons due to the theorem of Ref. 25. Letting M be the mass of a gauge boson of a broken symmetry and $\sqrt{s}$ the center-of-mass energy of a given process, then the theorem states that, to leading order in M^2/s, the S matrix for a process involving the longitudinal gauge boson is equivalent to that of the process with the longitudinal gauge boson replaced by its corresponding Goldstone boson.

The effective Lagrangian is constructed to be invariant under the $SU(3)_C\times SU(2)_L\times U(1)_Y$ gauge as well as the dilatation symmetries. We first realize the global SU(2)$\times$SU(2) chiral symmetry of the technicolor sector via the Goldstone bosons S^i which transform as

$$\delta_\chi S^i = A^i_\chi(S) \ . \tag{A1}$$

Here δ_χ denotes the SU(2)$\times$SU(2) variations and the Killing vectors $A^i_\chi(S)$ are given by

$$A^i_\chi(S)=\begin{cases}\epsilon_{aij}S^j, & \chi=a \text{ for } \delta_a \text{ a } SU(2)_V \text{ variation} ,\\ \delta^{ij}+S^iS^j, & \chi=j \text{ for } \delta_j \text{ a } SU(2)\times SU(2)/SU(2)_V \text{ variation} .\end{cases} \tag{A2}$$

Alternate choices of Killing vectors merely correspond to a reparametrization of the fields and leave the S matrix unchanged. The electroweak interactions are obtained by gauging the $SU(2)_L\times U(1)_Y$ subgroup of the global SU(2)$\times$SU(2) chiral symmetry. The Goldstone bosons corresponding to the physical $W^\pm$ and Z^0 bosons are defined as

$$S^\pm=\frac{1}{\sqrt{2}}(S^1\pm iS^2),\quad S^0=S^3 \ . \tag{A3}$$

To realize the scale transformation δ_D, the dilaton is defined to transform as the exponential field

$$U=e^{D/F_D} \ . \tag{A4}$$

The exponential scales as a canonical dimension one field

$$\delta_D U=(1+x^\lambda\partial_\lambda)U \tag{A5}$$

yielding the dilaton transformation equation

$$\delta_D D=F_D+x^\lambda\partial_\lambda D \ . \tag{A6}$$

The dilaton is, of course, a SU(2)$\times$SU(2) singlet. Since the scale and chiral transformations commute, we find that the chiral Goldstone bosons necessarily scale as dimension-zero fields

$$\delta_D S^i = x^\lambda\partial_\lambda S^i \ . \tag{A7}$$

All of the remaining fields transform as usual under the $SU(3)_C\times SU(2)_L\times U(1)_Y$ and technicolor gauge variations and with canonical scale dimensions under dilatation transformations.

As far as scale invariance is concerned, the effective Lagrangian is constructed from scale dimension four, $SU(3)_C\times SU(2)_L\times U(1)_Y$ gauge-invariant terms. Hence, none of the renormalizable fermion kinetic energy terms of the standard model will involve dilaton couplings.

The fermion mass terms are assumed to arise from the scale-invariant extended technicolor gauge coupling of the quarks and leptons to the technifermions. In turn we will assume that the effective-Lagrangian mass terms are scale invariant. Since the mass terms are dimension three they couple to one power of $U=e^{D/F_D}$. Their contribution to the effective action is

$$\mathscr{L}_m = e^{D/F_D} \sum_f \bar{f} m_f f \ . \tag{A8}$$

On the other hand, the running of the gauge coupling constants explicitly breaks the scale invariance of the theory. The dilaton couples directly to the breaking terms so that the effective Lagrangian reproduces the trace anomaly upon variation. The scale-breaking contribution to the effective Lagrangian is

$$\mathscr{L}_A = \frac{D}{F_D}\left[-\frac{1}{4}\frac{\beta_{\rm TC}^{\rm TF}}{g_{\rm TC}} T_{\mu\nu}T^{\mu\nu} - \frac{1}{4}\frac{\beta_2^{\rm TF}}{g_2} W_{\mu\nu}W^{\mu\nu} - \frac{1}{4}\frac{\beta_1^{\rm TF}}{g_1} B_{\mu\nu}B^{\mu\nu} \right] , \tag{A9}$$

where $\beta_{\rm TC}^{\rm TF}, \beta_2^{\rm TF}, \beta_1^{\rm TF}$ and $T_{\mu\nu}, W_{\mu\nu}, B_{\mu\nu}$ are, respectively, the renormalization-group β functions and field strengths of the technicolor, $\mathrm{SU(2)}_L$ and $\mathrm{U(1)}_Y$ gauge groups. The superscript TF denotes the fact that the β functions include only contributions from the technifermions.

The Goldstone-boson self-interactions involve two sets of terms. The scale dimension-four dilaton kinetic energy term is made from the U field and is given by

$$\mathscr{L}_D = \frac{{F_D}^2}{2}\partial_\mu U \partial^\mu U = \tfrac{1}{2} e^{2D/F_D} \partial_\mu D \partial^\mu D \ , \tag{A10}$$

while the chiral-symmetry Goldstone-boson $\mathrm{SU(2)}_L \times \mathrm{U(1)}_Y$ gauge-invariant action is[26]

$$\mathscr{L}_\chi = \tfrac{1}{2}(\partial_\mu S^i - V_{\mu\chi_1} A^i_{\chi_1}) g_{ij}(S)(\partial^\mu S^j - V^\mu_{\chi_2} A^j_{\chi_2}) \ . \tag{A11}$$

Here $g^{ij}(S)=A^i_\chi A^j_\chi$ is the metric for the $\mathrm{SU(2)}\times\mathrm{SU(2)}/\mathrm{SU(2)}_V$ coset space and the gauge bosons are defined as

$$V^\mu_\chi = \begin{cases} H^\mu_a, & \chi = a \in \mathrm{SU(2)}_V \ , \\ X^\mu_j, & \chi = j \in \mathrm{SU(2)}\times\mathrm{SU(2)}/\mathrm{SU(2)}_V \ . \end{cases} \tag{A12}$$

Rewriting in terms of the photon A, $W^\pm$, and Z^0 fields these take the form

$$\begin{aligned} X^\mu_1 &= \frac{-e}{2\sqrt{2}\sin\theta_w}(W^{+\mu}+W^{-\mu}) = -H^\mu_1 \ , \\ X^\mu_2 &= \frac{-ie}{2\sqrt{2}\sin\theta_w}(W^{+\mu}-W^{-\mu}) = -H^\mu_2 \ , \\ X^\mu_3 &= \frac{-e}{\sin 2\theta_w} Z^{0\mu} \ , \\ H^\mu_3 &= e(A^\mu + \cot 2\theta_w Z^{0\mu}) \ . \end{aligned} \tag{A13}$$

Since the chiral Goldstone bosons have scale dimension zero, $\mathscr{L}_\chi$ has only scale dimension two. Hence it couples to U^2 to yield the interactions of the dilaton with the electroweak gauge bosons which are given by

$$\begin{aligned} \mathscr{L}_{\rm GD} &= e^{2D/F_D}\mathscr{L}_\chi \ , \\ \mathscr{L}_\chi &= \frac{F^2}{F^2+S^2}(\partial^\mu S^+ \partial_\mu S^- + \tfrac{1}{2}\partial^\mu S^0 \partial_\mu S^0) + {M_W}^2 W^{+\mu}W^-_\mu + \tfrac{1}{2}{M_Z}^2 Z^{0\mu}Z^0_\mu + M_W(\partial^\mu S^+ W^-_\mu + \partial^\mu S^- W^+_\mu) \\ &\quad + M_Z \partial^\mu S^0 Z^0_\mu - \frac{1}{8}\frac{1}{F^2+S^2}(\partial^\mu S^2)(\partial_\mu S^2) + \frac{{M_W}^2}{F^2} S^2 W^{+\mu}W^-_\mu - i\frac{M_W}{F}[W^{+\mu}(S^- \overleftrightarrow{\partial}_\mu S^0) - W^{-\mu}(S^+ \overleftrightarrow{\partial}_\mu S^0)] \\ &\quad + i\frac{eM_W}{F}(A^\mu - \tan\theta_w Z^{0\mu})[F(S^- W^+_\mu - S^+ W^-_\mu) + iS^0(S^- W^+_\mu + S^+ W^-_\mu)] \\ &\quad + ie(A^\mu + \cot 2\theta_w Z^{0\mu})\{S^-[\overleftrightarrow{\partial}_\mu - ie(A_\mu + \cot 2\theta_w Z^0_\mu)]S^+\} + \frac{1}{2}\frac{{M_Z}^2}{F^2} S^{02} Z^{0\mu} Z^0_\mu \\ &\quad + \frac{M_W}{F^2}\left[S^+ W^{-\mu} + S^- W^{+\mu} + \frac{M_Z}{M_W} Z^{0\mu} \right](\partial_\mu S^2) \ , \end{aligned} \tag{A14}$$

with $S^2 = 2S^+S^- + S^{02}$.

The total $\mathrm{SU(3)}_C \times \mathrm{SU(2)}_L \times \mathrm{U(1)}_Y$ scale-invariant effective Lagrangian is thus

$$\mathscr{L}_{\rm eff} = \mathscr{L}_4 + \mathscr{L}_{\rm GD} + \mathscr{L}_D + \mathscr{L}_m + \mathscr{L}_A \ , \tag{A15}$$

where $\mathscr{L}_4$ are the dimension-four dilaton-independent standard-model and technigluon kinetic energy terms.

Since we are working in a renormalizable gauge we may also choose the gauge-fixing terms to be scale invariant. Typical R_ξ gauges will now have the chiral Goldstone field appear with products of U^2.

The scale-invariant R_ξ gauge-fixing terms in the effective action become

$$\mathscr{L}_\xi=\frac{1}{\xi}\left[\partial_\mu\left[\frac{W^{+\mu}F^2}{F^2+S^2}\right]+\xi M_W e^{2D/F_D}S^+\right]\left[\partial_\lambda\left[\frac{W^{-\lambda}F^2}{F^2+S^2}\right]+\xi M_W e^{2D/F_D}S^-\right]$$
$$+\frac{1}{2\xi_0}\left[\partial_\mu\left[\frac{Z^{0\mu}F^2}{F^2+S^2}\right]+\xi_0 M_Z e^{2D/F_D}S^0\right]^2+\frac{1}{2\alpha}(\partial_\mu A^\mu)^2 \ . \tag{A16}$$

In the standard manner this choice of gauge gives rise to Faddeev-Popov ghost and dilaton interaction terms which are irrelevant for our present discussion and hence will not be listed here.

---

[1]W. A. Bardeen, C. N. Leung, and S. T. Love, Phys. Rev. Lett. **56**, 1230 (1986).

[2]J. Kogut *et al.*, Phys. Rev. Lett. **48**, 1140 (1982); **50**, 393 (1983); Nucl. Phys. **B225**, 326 (1983); Phys. Lett. **145B**, 239 (1984); I. M. Barbour and C. J. Burden, *ibid.* **161B**, 357 (1985).

[3]For a discussion of this conjecture in the context of quenched, planar QED see C. N. Leung, S. T. Love, and W. A. Bardeen, Nucl. Phys. **B273**, 649 (1986).

[4]For studies of the dilaton in other contexts see, for example, R. J. Crewther, Phys. Lett. **33B**, 305 (1970); Phys. Rev. D **3**, 3152 (1971); J. Ellis, Nucl. Phys. **B22**, 478 (1970); A. A. Migdal and M. A. Shifman, Phys. Lett. **114B**, 445 (1982).

[5]S. Weinberg, Phys. Rev. D **13**, 974 (1976); **19**, 1277 (1979); L. Susskind, *ibid.* **20**, 2619 (1979).

[6]J. Ellis, M. K. Gaillard, and D. V. Nanopoulos, Nucl. Phys. **B106**, 292 (1976).

[7]A. D. Linde, Pis'ma Zh. Eksp. Teor. Fiz. **23**, 73 (1976) [JETP Lett. **23**, 64 (1976)]; S. Weinberg, Phys. Rev. Lett. **36**, 294 (1976).

[8]B. W. Lee, C. Quigg, and H. B. Thacker, Phys. Rev. Lett. **38**, 883 (1977); Phys. Rev. D **16**, 1519 (1977); M. Veltman, Acta Phys. Pol. **B8**, 475 (1977).

[9]B. Holdom, Phys. Rev. D **24**, 1441 (1981); K. Yamawaki, M. Bando, and K. Matumoto, Phys. Rev. Lett. **56**, 1335 (1986). See also T. Akiba and T. Yanagida, *ibid.* **169B**, 432 (1986).

[10]S. Dimopoulos and L. Susskind, Nucl. Phys. **B155**, 237 (1979); E. Eichten and K. Lane, Phys. Lett. **90B**, 125 (1980).

[11]See, however, T. W. Appelquist, D. Karabali, and L. C. R. Wijewardhana, Phys. Rev. Lett. **57**, 957 (1986).

[12]During the writing of the present paper we received a preprint by M. Bando, K. Matumoto, and K. Yamawaki, Report No. DPNU-86-13, 1986 (unpublished), in which the authors estimate the dilaton mass using the Dashen formula and obtain

$$m_D{}^2=-\frac{\beta(\alpha_{\rm TC})}{\alpha_{\rm TC}{}^2}\frac{\pi}{F_D{}^2}\left\langle\frac{\alpha_{\rm TC}}{\pi}F^{a\mu\nu}F_{a\mu\nu}\right\rangle_0=-\frac{\beta(\alpha_{\rm TC})}{\alpha_{\rm TC}{}^2}\Lambda_{\rm TC}{}^2 \ ,$$

where $F^{\mu\nu}$ denote the technigluon field strength. They then argue that since $\beta(\alpha_{\rm TC})/\alpha_{\rm TC}{}^2$ can be arbitrarily small near the nontrivial ultraviolet fixed point, $m_D \ll \Lambda_{\rm TC}$ and the dilaton should be very light. We disagree with this argument in that $\beta(\alpha_{\rm TC})/\alpha_{\rm TC}{}^2$ should be evaluated at $\mu^2=m_D{}^2$ which is far away from the ultraviolet limit. In fact, it may be closer to the infrared regime where $\beta(\alpha_{\rm TC})$ is unknown. Moreover, the above authors do not seem to appreciate the necessary presence of at least two separate scales ($\mu_{\rm TC}$ and $\Lambda_{\rm TC}$) in order for the dilaton to exist.

[13]The upper bound of 50 GeV is chosen so that m_D is below the $t\bar{t}$ threshold, assuming $m_t > 25$ GeV.

[14]D. Kohler, B. A. Watson, and J. A. Becker, Phys. Rev. Lett. **33**, 1628 (1974); S. J. Freedman, J. Napolitano, J. Camp, and M. Kroupa, *ibid.* **52**, 240 (1984); R. Barbieri and T. E. O. Ericson, Phys. Lett. **57B**, 270 (1975).

[15]A. I. Vainshtein, V. I. Zakharov, and M. A. Shifman, Usp. Fiz. Nauk **131**, 537 (1980) [Sov. Phys. Usp. **23**, 429 (1980)]; R. S. Willey and H. L. Yu, Phys. Rev. D **26**, 3287 (1982).

[16]T. N. Pham and D. G. Sutherland, Phys. Rev. D **34**, 1634 (1986); Phys. Lett. **151B**, 444 (1985).

[17]R. S. Willey, Phys. Lett. **173B**, 480 (1986).

[18]M. I. Vysotsky, Phys. Lett. **97B**, 159 (1980).

[19]This is anticipated since the coupling of the dilaton to light-quark states should be affected by the large explicit breaking of the scale symmetry at low energies which are neglected in the estimate of Eq. (10).

[20]Vainshtein, Zakharov, and Shifman (Ref. 15); A. A. Ansel'm, N. G. Ural'tsev, and V. A. Khoze, Usp. Fiz. Nauk. **145**, 185 (1985) [Sov. Phys. Usp. **28**, 113 (1985)].

[21]J. Ellis, K. Enqvist, D. V. Nanopoulos, and S. Ritz, Phys. Lett. **158B**, 417 (1985).

[22]R. N. Cahn, M. S. Chanowitz, and N. Fleishon, Phys. Lett. **82B**, 113 (1979).

[23]H. M. Georgi, S. L. Glashow, M. E. Machacek, and D. V. Nanopoulos, Phys. Rev. Lett. **40**, 692 (1978).

[24]M. S. Chanowitz and M. K. Gaillard, Phys. Lett. **142B**, 85 (1984).

[25]J. M. Cornwall, D. N. Levin, and G. Tiktopoulos, Phys. Rev. D **10**, 1145 (1974); B. W. Lee, C. Quigg, and H. B. Thacker, *ibid.* **16**, 1519 (1977).

[26]D. Boulware and L. S. Brown, Ann. Phys. (N.Y.) **138**, 392 (1982).