Collective physics in the closed bosonic string

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A second-quantized analysis is performed to examine many-body phenomena in closed bosonic strings. The covariant nonpolynomial closed-string field theory is developed in terms of particle fields and shown to contain interactions triggering a nonperturbative condensation of the tachyon field. We study the possibility that the higher-dimensional Lorentz symmetry spontaneously breaks. We show that the theory has asymptotic freedom due to a tree-level running coupling. The spectrum of states in the nonperturbative ground state is radically changed relative to the free case; in particular, there is no massless graviton. Similar effects are anticipated in any nonperturbative vacuum.

I. INTRODUCTION

The widespread interest in string theory[1] stems from remarkable features, present in the theory at the first-quantized level, that provide the possibility of combining gravity with other fundamental fields in a single self-consistent structure. However, a first-quantized investigation of any interacting theory is insufficient, since collective effects can substantially alter the physics. A well-known illustration is provided by the SU(2) scalar doublet of the standard model, which is tachyonic at the first-quantized level but which condenses in the field-theory vacuum due to self-interactions. The resulting physics involves no tachyons, while an effective mass is generated for the weak interactions.

The development of a covariant string field theory for the open bosonic string[2] made it possible to initiate a field-theoretic analysis[3] of collective open-string effects.[4] This semiclassical investigation indicates that many scalars, including the tachyon, condense in a nonperturbative open-string vacuum. The physics in the new vacuum is radically different from the first-quantized picture: only massive states occur, and certain states present in the canonical vacuum are absent.

The goal of the present work is to determine whether these features extend to the case of the closed bosonic string. This is feasible because a covariant closed-string field theory has recently been constructed.[5] It is nonpolynomial, i.e., it involves all powers of the string field. It is known to produce correct tree-level results, which suffices for a semiclassical study.

The existence of the closed-string tachyon destabilizes the canonical 26-dimensional vacuum. In analogy to the standard model, collective effects may stabilize the closed string in another ground state, containing a nonzero tachyon expectation. Such an expectation value is necessarily nonperturbative in the string coupling and the Regge slope.[6]

The results presented here are obtained through a mass-level truncation scheme. In the context of the open bosonic string, this scheme provides a systematic and self-consistent approach.[4] It is guided by the intuitive idea that light states dominate physical processes and is plausible in part because the particle-field couplings in covariant string-field theory are exponentially suppressed by the total level number of the fields involved. Additional evidence arises from analytical studies and from the convergence of numerical computations.

In the present context, the truncation scheme is systematic and self-consistent when all orders in the nonpolynomial interaction are included. It turns out that the couplings are further suppressed by the total level number relative to the open-string case. However, even at the lowest truncation level an all-orders calculation is impractical. Fortunately, qualitative features likely to be present in the full theory can be observed at low polynomial order. We limit discussion here to an explicit study of the cubic couplings in a truncation containing the tachyon and massless fields. This suffices for our purposes and avoids excessive complexity.

Section II presents the particle-field Lagrangian at this level and explores its vacuum structure. Its physics is analyzed in Sec. III. A summary and discussion is provided in Sec. IV.

II. A CANDIDATE NONPERTURBATIVE VACUUM

The action for the nonpolynomial closed-string field theory has the form[5]

$$S=\int\left[\frac{1}{\alpha'}\Phi * Q b_0^- \Phi + g\sum_{N=3}^{\infty}\frac{(g\alpha'/2)^{N-3}}{N!}\Phi * [\Phi^{N-1}]\right] . \tag{2.1}$$

It is invariant under the string-field gauge transformation

$$\delta(b_0^-\Phi)=\left[\frac{2}{\alpha'}\right]^{1/2} Q b_0^- \Lambda + \left[\frac{\alpha'}{2}\right]^{1/2} g\sum_{N=3}^{\infty}\frac{(g\alpha'/2)^{N-3}}{(N-2)!}[\Phi^{N-2}\Lambda] . \tag{2.2}$$

In these equations, Φ is the closed-string field, Q is the Becchi-Rouet-Stora-Tyutin (BRST) operator, $b_0^- = (b_0 - \bar{b}_0)/2$ is a combination of antighost zero modes, and Λ is the string gauge field. The on-shell three-tachyon coupling is denoted by g and the Regge slope is α'. The symbols $\int \Phi_1 * \Phi_2$ represent the string-field scalar product, and $[\Phi_1 \cdots \Phi_{N-1}]$ represents the string field obtained from the combination of the $N-1$ string fields $\Phi_1, \ldots, \Phi_{N-1}$ using the N-string vertex function.

At the chosen truncation level, the string field Φ has the following expansion in terms of particle fields:

$$\Phi = +c_0^- \left[\phi + A_{\mu\nu}\alpha^{\mu}_{-1}\bar{\alpha}^{\nu}_{-1} + \frac{1}{\sqrt{2}}(\alpha_+ + \alpha_-) b_{-1}\bar{c}_{-1} + \frac{1}{\sqrt{2}}(\alpha_+ - \alpha_-) c_{-1}\bar{b}_{-1} \right] |0\rangle$$
$$+ c_0^- c_0^+ (i j_{1\mu}\alpha^{\mu}_{-1}\bar{b}_{-1} + i j_{2\mu} b_{-1}\bar{\alpha}^{\mu}_{-1})|0\rangle \; . \tag{2.3}$$

Here, ϕ represents the tachyon field, $A_{\mu\nu}$ represents the two-tensor field, and $\alpha_\pm$, $j_{1\mu}$, and $j_{2\mu}$ represent auxiliary fields. The symmetric part of $A_{\mu\nu}$ is the graviton $h_{\mu\nu}$, which is a fluctuation about a flat background. The first-quantized string vacuum is $|0\rangle = \bar{c}_1 |\bar{\Omega}\rangle c_1 |\Omega\rangle$, where $|\bar{\Omega}\rangle$ and $|\Omega\rangle$ are the left and right Sl(2)-invariant vacua, respectively.[7] The coefficients of the fields are first-quantized oscillators; in particular, $c_0^- = c_0 - \bar{c}_0$ and $c_0^+ = (c_0 + \bar{c}_0)/2$.

To this truncation order, the quadratic Lagrangian in terms of particle fields is

$$\mathcal{L}_{\text{free}} = +\tfrac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{2}{\alpha'}\phi^2 + \tfrac{1}{2}\partial_\lambda A_{\mu\nu}\partial^\lambda A^{\mu\nu} + \tfrac{1}{2}\partial_\mu\alpha_+\partial^\mu\alpha_+ - \tfrac{1}{2}\partial_\mu\alpha_-\partial^\mu\alpha_-$$
$$+ \frac{1}{\sqrt{2\alpha'}}(j_{1\mu}\partial_\nu A^{\mu\nu} + j_{2\nu}\partial_\mu A^{\mu\nu}) + \frac{1}{\alpha'}(j_{1\mu}j_1^\mu + j_{2\mu}j_2^\mu) \; . \tag{2.4}$$

The cubic Lagrangian is lengthy. To simplify matters, we present it in the Siegel-Feynman gauge[8] $b_0^+\Phi = 0$; this choice sets the particle fields $j_{1\mu}$ and $j_{2\mu}$ to zero. Even with this simplification, the cubic Lagrangian contains about 50 terms at this level. We present these terms arranged according to the total level number, where the tachyon is at level zero and the other fields are at level two:

$$\mathcal{L}_{\text{cubic}} = \mathcal{L}^{(0)} + \mathcal{L}^{(2)} + \mathcal{L}^{(4)} + \mathcal{L}^{(6)} \; . \tag{2.5}$$

We find

$$\mathcal{L}^{(0)} = \frac{1}{3!}\epsilon^{-3} g\tilde{\phi}^3 \; , \tag{2.6}$$

where the suppression factor $\epsilon = 2^4/3^3 \approx 0.6$. For the order-two terms we have

$$\mathcal{L}^{(2)} = \frac{1}{2^3}\epsilon^{-2} g\alpha'(\partial_\mu\tilde{\phi}\partial_\nu\tilde{\phi}\tilde{A}^{\mu\nu} - \tilde{\phi}\partial_\mu\partial_\nu\tilde{\phi}\tilde{A}^{\mu\nu}) \; . \tag{2.7}$$

The order-four terms are

$$\mathcal{L}^{(4)} = +\tfrac{1}{2}\epsilon^{-1} g\tilde{\phi}\tilde{A}_{\mu\nu}\tilde{A}^{\mu\nu} + \frac{1}{2^3}\epsilon^{-1} g\tilde{\phi}(\tilde{\alpha}_+^2 - \tilde{\alpha}_-^2)$$
$$+ \frac{1}{2^4}\epsilon^{-1} g\alpha'(\tilde{\phi}\partial_\mu\tilde{A}_{\nu\lambda}\partial^\nu\tilde{A}^{\mu\lambda} - 2\partial_\mu\tilde{\phi}\tilde{A}_{\nu\lambda}\partial^\nu\tilde{A}^{\mu\lambda} + \partial_\mu\partial_\nu\tilde{\phi}\tilde{A}^{\mu}{}_{\lambda}\tilde{A}^{\nu\lambda}$$
$$+ \tilde{\phi}\partial_\mu\tilde{A}_{\lambda\nu}\partial^\nu\tilde{A}^{\lambda\mu} - 2\partial_\mu\tilde{\phi}\tilde{A}_{\lambda\nu}\partial^\nu\tilde{A}^{\lambda\mu} + \partial_\mu\partial_\nu\tilde{\phi}\tilde{A}_{\lambda}{}^{\mu}\tilde{A}^{\lambda\nu})$$
$$+ \frac{1}{2^7}\epsilon^{-1} g\alpha'^2(\tilde{\phi}\partial_\mu\partial_\nu\tilde{A}_{\rho\sigma}\partial^\rho\partial^\sigma\tilde{A}^{\mu\nu} - 2\partial_\mu\tilde{\phi}\partial_\nu\tilde{A}_{\rho\sigma}\partial^\rho\partial^\sigma\tilde{A}^{(\mu\nu)} + 2\partial_\mu\partial_\nu\tilde{\phi}\tilde{A}_{\rho\sigma}\partial^\rho\partial^\sigma\tilde{A}^{\mu\nu}$$
$$+ \partial_\mu\partial_\nu\tilde{\phi}\partial_\rho\tilde{A}^{(\mu\sigma)}\partial_\sigma\tilde{A}^{(\nu\rho)} - 2\partial_\mu\partial_\nu\partial_\sigma\tilde{\phi}\tilde{A}^{(\rho\sigma)}\partial_\rho\tilde{A}^{\mu\nu} + \partial_\mu\partial_\nu\partial_\rho\partial_\sigma\tilde{\phi}\tilde{A}^{\mu\nu}\tilde{A}^{\rho\sigma}) \; . \tag{2.8}$$

Parentheses enclosing indices denote symmetrization, e.g., $\tilde{A}^{(\mu\nu)} = \tilde{A}^{\mu\nu} + \tilde{A}^{\nu\mu} = h^{\mu\nu}$. Finally, for the order-six terms we have

$$\mathcal{L}^{(6)} = -\frac{1}{2^5} g\alpha'\tilde{A}_{\mu\nu}(\tilde{\alpha}_+\partial^\mu\partial^\nu\tilde{\alpha}_+ - \tilde{\alpha}_-\partial^\mu\partial^\nu\tilde{\alpha}_- - \partial^\mu\tilde{\alpha}_+\partial^\nu\tilde{\alpha}_+ + \partial^\mu\tilde{\alpha}_-\partial^\nu\tilde{\alpha}_-) - \frac{1}{2^3} g\alpha'\tilde{A}_{\mu\nu}(\tilde{A}_{\rho\sigma}\partial^\mu\partial^\nu\tilde{A}^{\rho\sigma} - \partial^\mu\tilde{A}_{\rho\sigma}\partial^\nu\tilde{A}^{\rho\sigma})$$
$$+ \frac{1}{2^3} g\alpha'\tilde{A}_{\mu\nu}(\tilde{A}_{\rho\sigma}\partial^\mu\partial^\sigma\tilde{A}^{\rho\nu} + \partial^\rho\tilde{A}^{\mu\sigma}\partial_\sigma\tilde{A}_{\rho\nu} - \partial^\rho\tilde{A}^{\mu\sigma}\partial^\nu\tilde{A}_{\rho\sigma} - \partial^\mu\tilde{A}_{\rho\sigma}\partial^\sigma\tilde{A}^{\rho\nu})$$
$$+ \frac{1}{2^6} g\alpha'^2(\partial_\tau\tilde{A}_{\mu\nu}\partial^\mu\partial^\nu\tilde{A}_{\rho\sigma}\partial^\sigma\tilde{A}^{\rho\tau} - \partial^\sigma\partial_\tau\tilde{A}_{\mu\nu}\partial^\mu\partial^\nu\tilde{A}_{\rho\sigma}\tilde{A}^{\rho\tau} - \partial_\tau\tilde{A}_{\mu\nu}\partial^\mu\tilde{A}_{\rho\sigma}\partial^\nu\partial^\sigma\tilde{A}^{\rho\tau}$$
$$+ \partial^\sigma\partial_\tau\tilde{A}_{\mu\nu}\partial^\mu\tilde{A}_{\rho\sigma}\partial^\nu\tilde{A}^{\rho\tau} - \partial_\tau\tilde{A}_{\mu\nu}\partial^\nu\tilde{A}_{\rho\sigma}\partial^\mu\partial^\sigma\tilde{A}^{\rho\tau} + \partial^\tau\tilde{A}_{\mu\nu}\tilde{A}_{\rho\sigma}\partial^\mu\partial^\nu\partial^\sigma\tilde{A}^{\rho\tau}$$
$$- \tilde{A}_{\mu\nu}\partial^\mu\partial^\nu\partial^\tau\tilde{A}_{\rho\sigma}\partial^\sigma\tilde{A}^{\rho\tau} + \tilde{A}_{\mu\nu}\partial^\mu\partial_\tau\tilde{A}_{\rho\sigma}\partial^\nu\partial^\sigma\tilde{A}^{\rho\tau} + \partial_\tau\tilde{A}_{\mu\nu}\partial^\mu\partial^\nu\tilde{A}_{\sigma\rho}\partial^\sigma\tilde{A}^{\tau\rho} - \partial^\sigma\partial_\tau\tilde{A}_{\mu\nu}\partial^\mu\partial^\nu\tilde{A}_{\sigma\rho}\tilde{A}^{\tau\rho}$$
$$- \partial_\tau\tilde{A}_{\mu\nu}\partial^\mu\tilde{A}_{\sigma\rho}\partial^\nu\partial^\sigma\tilde{A}^{\tau\rho} + \partial^\sigma\partial_\tau\tilde{A}_{\mu\nu}\partial^\mu\tilde{A}_{\sigma\rho}\partial^\nu\tilde{A}^{\tau\rho} - \partial_\tau\tilde{A}_{\mu\nu}\partial^\nu\tilde{A}_{\sigma\rho}\partial^\mu\partial^\sigma\tilde{A}^{\tau\rho} + \partial^\tau\tilde{A}_{\mu\nu}\tilde{A}_{\sigma\rho}\partial^\mu\partial^\nu\partial^\sigma\tilde{A}^{\tau\rho}$$
$$- \tilde{A}_{\mu\nu}\partial^\mu\partial^\nu\partial^\tau\tilde{A}_{\sigma\rho}\partial^\sigma\tilde{A}^{\tau\rho} + \tilde{A}_{\mu\nu}\partial^\mu\partial_\tau\tilde{A}_{\sigma\rho}\partial^\nu\partial^\sigma\tilde{A}^{\tau\rho})$$

$$-\frac{1}{2^{10}3}g\alpha'^3(6\tilde{A}_{\lambda\mu}\partial^\sigma\partial^\tau\tilde{A}_{\nu\rho}\partial^\lambda\partial^\mu\partial^\nu\partial^\rho\tilde{A}_{\sigma\tau}-6\tilde{A}_{\lambda\mu}\partial^\mu\partial^\sigma\partial^\tau\tilde{A}_{\nu\rho}\partial^\lambda\partial^\nu\partial^\rho\tilde{A}_{\sigma\tau}-6\partial^\sigma\tilde{A}_{\lambda\mu}\partial^\tau\tilde{A}_{\nu\rho}\partial^\lambda\partial^\mu\partial^\nu\partial^\rho\tilde{A}_{\sigma\tau}$$
$$-6\partial^\rho\tilde{A}_{\lambda\mu}\partial^\sigma\partial^\tau\tilde{A}_{(\nu\rho)}\partial^\lambda\partial^\mu\partial^\nu\tilde{A}_{\sigma\tau}+6\partial^\sigma\tilde{A}_{(\lambda\mu)}\partial^\mu\partial^\tau\tilde{A}_{\nu\rho}\partial^\lambda\partial^\nu\partial^\rho\tilde{A}_{(\sigma\tau)}$$
$$+2\partial^\sigma\partial^\tau\tilde{A}_{\lambda\mu}\partial^\lambda\partial^\mu\tilde{A}_{\nu\rho}\partial^\nu\partial^\rho\tilde{A}_{\sigma\tau}-\partial^\rho\partial^\tau\tilde{A}_{(\lambda\mu)}\partial^\lambda\partial^\sigma\tilde{A}_{(\nu\rho)}\partial^\mu\partial^\nu\tilde{A}_{(\sigma\tau)})\ . \quad (2.9)$$

In Eq. (2.9), the terms cubic in $A_{\mu\nu}$ and linear in α' contain the standard Einstein gravitational interactions in the Feynman-Siegel gauge; other terms represent string corrections. Note that all fields f entering the interaction Lagrangian are smeared over a distance $\sqrt{\alpha'}$:

$$\tilde{f}=\exp[\tfrac{1}{2}\alpha'\ln(3\sqrt{3}/4)\partial_\mu\partial^\mu]f\ . \quad (2.10)$$

We seek ground states of the theory, which are suitable extrema of the static potential. The latter consists of all momentum-independent terms in the Lagrangian. We find

$$V_{\text{static}}=-\frac{2}{\alpha'}\phi^2+\frac{1}{3!}\epsilon^{-3}g\phi^3+\frac{1}{2^3}\epsilon^{-1}g\phi(2^2A_{\mu\nu}A^{\mu\nu}+\alpha_+^2-\alpha_-^2)\ . \quad (2.11)$$

The canonical 26-dimensional vacuum is an unstable extremum of V_{static}: $\langle\phi\rangle=\langle A_{\mu\nu}\rangle=\langle\alpha_\pm\rangle=0$. However, a local minimum occurs at $\langle\phi\rangle=2^{15}/3^9g\alpha'\approx1.66/g\alpha'$, $\langle A_{\mu\nu}\rangle=0,\langle\alpha_\pm\rangle=0$. It can be shown to be perturbatively stable. The new ground state is nonperturbative because $\langle\phi\rangle$ is of order $1/g$. It generates a contribution $\approx-4.77/(g^2\alpha'^3)$ to the cosmological constant.

The existence of this vacuum is unaffected by the gauge choice. If no gauge condition is imposed, the static potential is augmented by the terms

$$\mathcal{L}_{\text{static}}^{\text{gauge}}=-\frac{1}{2^3}\epsilon^{-1}g\phi(j_{1\mu}j_1^\mu+j_{2\mu}j_2^\mu)-\frac{1}{2^2}gA^{\mu\nu}j_{1\mu}j_{2\nu}+\frac{1}{2^3\sqrt{2}}g[(\alpha_++\alpha_-)j_{1\mu}j_1^\mu-(\alpha_+-\alpha_-)j_{2\mu}j_2^\mu)]\ . \quad (2.12)$$

The candidate nonperturbative vacuum found above is now specified also by the conditions $j_{1\mu}=j_{2\mu}=0$. Other nonperturbative extrema exist, but they are either unstable or are gauge artifacts.

III. COLLECTIVE PHYSICS

The interaction Lagrangian $\mathcal{L}_{\text{cubic}}$ includes terms that are purely stringy in the sense that they are forbidden by symmetries in an ordinary particle-gravity theory. For example, the term $\phi A_{\mu\nu}A^{\mu\nu}$ appearing in Eq. (2.11) incorporates a tachyon-graviton-graviton piece that would normally be incompatible with general coordinate invariance. Terms of this type can coexist with symmetries in string theories because invariance is achieved through cancellations among contributions from infinitely many interaction terms.

In principle, such scalar-tensor-tensor terms provide a natural means for the spontaneous breaking of the 26-dimensional spacetime symmetry,[9] which might be a signal of spontaneous compactification and could result in phenomena observable in solar-system experiments. To illustrate, if ϕ acquires a negative expectation $\langle\phi\rangle$ then the tachyon-graviton-graviton term $\phi A_{\mu\nu}A^{\mu\nu}$ appearing in Eq. (2.11) generates a negative squared mass for $A_{\mu\nu}$. The end result is a nonzero graviton condensate $\langle A_{\mu\nu}\rangle$, which spontaneously breaks the higher-dimensional Lorentz symmetry. In fact, this effect does not occur for the candidate ground state presented in Sec. II because $\langle\phi\rangle$ is positive.

Another stringy feature of the interaction Lagrangian, valid also in the full nonpolynomial theory, is the smearing of particle fields over a distance of order of the Planck length seen in Eq. (2.10). It follows that the momentum-space couplings carry explicit momentum dependence. This means that the effective coupling $g(p^2)$ runs already at tree level in a string theory. Indeed, the exponential decrease of $g(p^2)$ for large spatial Euclidean p^2 or, equivalently, for small spatial distances on the Planck scale makes the closed bosonic string asymptotically free.

The momentum dependence of the effective coupling has several implications for the theory. For one, the determination of tree-level scattering processes in any nonperturbative vacuum requires both mass and wave-function renormalization. Another is that phenomena at high p^2 in any vacuum are perturbatively calculable. Still another feature of the running coupling is that the effective coupling becomes exponentially strong for sufficiently large timelike momenta. This means that perturbation theory cannot be used to determine the spectrum of states at high-mass levels.

A particularly relevant consequence of the momentum dependence of $g(p^2)$ is its effect on the mass spectrum in *all* coupling regimes. The asymptotic Hilbert space in a nonperturbative vacuum is radically affected because most propagators acquire a transcendental structure, which eliminates poles for certain states.

A simple example is provided by the tachyon propagator in the level-zero truncation. Shifting by the tachyon expectation and collecting quadratic terms yields the inverse tachyon propagator in Euclidean momentum space as

$$Q_\phi(p^2)=\frac{1}{2\alpha'}\{\alpha'p^2-2^2+2^3\exp[-\alpha'p^2\ln(3\sqrt{3}/4)]\}\ . \quad (3.1)$$

Without the exponential smearing, there would be a Minkowski-space pole at $m^2=4/\alpha'$. However, an explicit examination of Eq. (3.1) shows that the transcendental behavior precludes a contribution to the spectrum of asymptotic states.[10]

It can be seen from the Lagrangian presented in Sec. II that the tachyon field mixes with fields at mass-level zero when the latter are included. Mixing of fields from different mass levels is generic in nonperturbative vacua. Note, however, that tachyon mixing with mass-level-zero fields does *not* occur in the open bosonic string.[4]

As another example with more physical import, consider the transverse traceless part $\hat{h}^{TT}_{\mu\nu}$ of $h_{\mu\nu}$ given by

$$\hat{h}^{TT}_{\mu\nu}=\hat{h}_{\mu\nu}-\frac{\partial_{(\mu}\partial^{\lambda}\hat{h}_{\lambda\nu)}}{\partial^2}+\frac{1}{25}\left[\eta_{\mu\nu}+24\frac{\partial_\mu\partial_\nu}{\partial^2}\right]\frac{\partial^\rho\partial^\sigma}{\partial^2}\hat{h}_{\rho\sigma}\ ,$$
$$\hat{h}_{\mu\nu}=h_{\mu\nu}-\tfrac{1}{26}\eta_{\mu\nu}\eta_{\rho\sigma}h^{\rho\sigma}\ . \tag{3.2}$$

This component carries the physical gravitational interaction. Shifting into the nonperturbative vacuum and collecting quadratic terms in $\hat{h}^{TT}_{\mu\nu}$ from the full Lagrangian in Sec. II generates the quadratic form $Q_h(p^2)$ whose inverse is the physical graviton propagator.

$$Q_h(p^2)=\frac{1}{2\alpha'}\left[\alpha'p^2+\frac{2^{11}}{3^6}\exp[-\alpha'p^2\ln(3\sqrt{3}/4)]\right]\ . \tag{3.3}$$

A direct check demonstrates that no asymptotic graviton pole occurs.[10] The graviton apparently cannot propagate beyond the Planck scale in the nonperturbative vacuum.

We conclude this section with two remarks. First, the results presented above suggest that the spectrum of the full nonperturbative closed-string field theory in a nonperturbative vacuum may contain no tachyons and no massless states. If true, this means that there are no infrared divergences and the theory is most likely finite. Second, a detailed analysis of the propagators in the nonperturbative vacuum indicates that it is perturbatively stable. However, the possibility exists that tunneling to another vacuum may occur, so that the nonperturbative vacuum may be nonperturbatively unstable. It is likely that penetration of a barrier is heavily suppressed, if it occurs at all, because it must involve the entire string, i.e., an infinite number of particle fields with a penetration-suppression factor for each.

IV. DISCUSSION

In this work, we have demonstrated that stringy many-body effects arise in the second-quantized analysis of the closed bosonic string. The string self-interactions induce the formation of a nonperturbative ground state in which the tachyon acquires a nonzero expectation value. The physics seen in the new vacuum is substantially different from that in the canonical one. Among the effects are the running of the tree-level string coupling. This causes asymptotic freedom and has consequences due to strong coupling for the perturbative analysis of the high-mass spectrum in any vacuum, including the canonical one, once interactions are present. Even in a weak-coupling regime, the Hilbert space is substantially changed. Propagator poles become renormalized and can disappear. An important example is the graviton field, which has no massless pole in the new vacuum.

Many of these effects are expected to occur in nonperturbative vacua of any string theory. Indeed, similar results have been obtained in a study of collective effects in the open bosonic string.[4] Since the exponential running of the string coupling is an immediate consequence of the extended nature of the string, the Hilbert space of any interacting string is likely to be different from the free limit. For the same reason, asymptotic freedom should also be generic.

The full nonpolynomial field theory of the closed bosonic string contains many stringy couplings involving powers of scalar fields with quadratic terms in the graviton. As we have seen, these couplings are stringy in the sense that they are excluded by general coordinate invariance in standard particle-gravity models. Since the canonical closed-string vacuum is unstable, nonperturbative scalar condensates form. The graviton is affected in one of three ways.[11] First, the graviton itself can form a condensate. This spontaneously breaks the 26-dimensional Lorentz symmetry. Second, the graviton may acquire a mass. The natural scale for this mass is the Planck scale, so barring fortuitous cancellations gravity is a short-range force in this scenario. Third, it may be that no graviton pole appears in the asymptotic spectrum. The latter two situations are undesirable if the closed string is to be a theory of quantum gravity. Evidently, the possibility of collective effects of the type presented means that general coordinate invariance is insufficient to guarantee a massless graviton in string theory.

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[1]See, for example, M. B. Green, J. H. Schwarz, and E. Witten, *Superstring Theory* (Cambridge University Press, Cambridge, England, 1987), Vols. I and II.

[2]E. Witten, Nucl. Phys. **B268**, 253 (1986).

[3]Some explorations of collective effects have been performed in the context of *first*-quantized physics. Among these are the ingenious ideas of M. B. Green, Phys. Lett. B **201**, 42 (1988); Nucl. Phys. **B103**, 333 (1976); **B124**, 461 (1977), involving a change from Neumann to Dirichlet boundary conditions on the world sheet; and the works of K. Bardakçi, Nucl. Phys. **B68**, 331 (1974); **B70**, 397 (1974); **B133**, 297 (1978), and of K. Bardakçi and M. B. Halpern, Phys. Rev. D **10**, 4230 (1974); Nucl. Phys. **B73**, 295 (1974); **B96**, 285 (1975).

[4]V. A. Kostelecký and S. Samuel, Phys. Rev. Lett. **63**, 224 (1990); Nucl. Phys. **B336**, 263 (1990).

[5]T. Kugo and K. Suehiro, Nucl. Phys. **B337**, 434 (1990); T. Kugo, H. Kunimoto, and K. Suehiro, Phys. Lett. B **226**, 48 (1989); M. Saadi and B. Zweibach, Ann. Phys. (N. Y.) **192**, 213 (1989); M. Kaku, in *Functional Integration, Geometry and Strings*, edited by Z. Haba and J. Sobczyk (Birkhauser, Boston, 1989); Phys. Rev. D **41**, 3734 (1990).

[6]V. A. Kostelecký and S. Samuel, Phys. Lett. B **207**, 169 (1988).

[7]D. Friedan, E. Martinec, and S. Shenker, Nucl. Phys. **271**, 93 (1986).

[8]W. Siegel, Phys. Lett. **142B**, 276 (1984); **151B**, 391 (1985); **151B**, 396 (1985).

[9]V. A. Kostelecký and S. Samuel, Phys. Rev. D **39**, 683 (1989); Phys. Rev. Lett. **63**, 224 (1989); Phys. Rev. D **40**, 1886 (1989).

[10]In principle, the quadratic form might have complex poles entering a physical *S*-matrix element. These could be manifested as inconsequential enhancements of the background in scattering or as the minute causality violations that might be expected from any nonlocal, relativistic theory.

[11]One of these three possibilities must also apply to the antisymmetric two-tensor field.