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STRING FIELD THEORY IN COMPONENTS: GENERAL COVARIANCE
VERSUS MASSIVE FIELDS

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String field theory in components:
general covariance versus massive fields

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Abstract

Arguments are given for non-existence of a field theory action for the massless and infinite (or a finite but non-zero) number of massive modes which satisfies the two requirements: (i) it reproduces closed string-theory S-matrix both in the massless and the massive sectors; (ii) it is invariant under the standard general coordinate transformations with the massive fields transforming simply as tensors. Connections between string field theory, path integral and σ -model approaches to derivation of string equations of motion are discussed.

1. In this paper we shall discuss whether it is possible to represent (at the tree level) a string theory by a generally covariant theory of infinite number of local fields corresponding to string excitation modes. The basic requirement on the field theory is that it must reproduce a (tree-level) string-theory S-matrix. The standard dual model picture suggests that tree string diagrams can be represented as trees in an infinite-component Φ^3 -type field theory [1]. We shall consider only the closed Bose string and assume that a Φ^3 -type string field theory action exists in this case [1-6]. Near a trivial flat-space vacuum we have: $S = \frac{1}{g^2} \int (\Phi_\Delta \Phi + \Phi^3)$. Representing the string field Φ in terms of a finite number of massless (A) and an infinite number of massive (B) fields using the standard normal mode expansion we get, symbolically,

$$S = \frac{1}{g^2} \int d^2x \left\{ A \square A + B(\square + m^2)B + A^3 + A^2 B + A B^2 + B^3 \right\} \quad (1)$$

For simplicity we consider A and B as two scalar fields and ignore complicated derivative structure of the vertices (which are essentially 3-point string amplitudes extended off-shell). While A includes graviton $h_{\mu\nu}$, (1) is not invariant under the ordinary general coordinate transformations (e.g. it lacks the Einstein term). At the same time, the string field theory is invariant under generalized gauge transformations which include the linearized general coordinate transformations at the linearized level but mix A's and B's at the non-linear level. The non-linear general covariance appears only as an approximate symmetry (present at energies $\ll m \sim 1/\sqrt{\alpha'}$) after we eliminate ("integrate over") all massive fields from (1) and from the gauge transformation laws using their equations of motion. Let $B_0(A) = -(\square + m^2)^{-1} A^2 + \dots$ be a solution of $\frac{\delta S}{\delta B} = 0$. If we make a change of massive variables, $B = \tilde{B} + B_0(A)$, then the action (1) takes the form (see also [7])

$$S = S_{\text{cov}}[A] + S_2[\tilde{B}, A] + S_3[\tilde{B}] \quad , \quad (2)$$

$$S_{\text{cov}} = g^{-2} \int d^3x \left\{ A \square A + A^3 - A^2 (\square + m^2)^{-1} A^2 + \dots \right\}, \quad (3)$$

$$S_2 = g^{-2} \int d^3x \tilde{B} \{ \square + m^2 + A + B_0(A) \} \tilde{B}, \quad S_3 = g^{-2} \int d^3x \tilde{B}^3. \quad (4)$$

Here S_{cov} (expanded in powers of $\mathcal{A}'$) is a covariant effective action for the massless fields [8]. The transformation $B \rightarrow \tilde{B}$ eliminated the $A^2 B$ -coupling in S and thus "decoupled" the massless sector from the massive. While S_2 (containing arbitrary powers of A) could in principle turn out to be covariant under the standard coordinate transformations (with $\tilde{B}$'s transforming simply as tensors), S_3 is certainly non-covariant:

$$\tilde{B}^3 \sim \sum \tilde{B}_{\dots} \partial \dots \partial \tilde{B}_{\dots} \partial \dots \partial \tilde{B}_{\dots}$$

depends on the flat (and not the full $G_{\mu\nu} = \delta_{\mu\nu} + h_{\mu\nu}$) metric.

While (2) correctly reproduces the string S -matrix in the A -sector it fails (in contrast to (1)) to do so in the sectors with $\tilde{B}$ -legs. The reason is that the transformation $B = \tilde{B} + B_0(A)$, being non-local, spoils the conditions of the "equivalence theorem" [9] (the S -matrices corresponding to (1) and (2) differ already at 3-point level since (2) lacks $A^2 \tilde{B}$ -vertex).

Let us now give a simple argument in favor of non-existence of a generally covariant field theory action which reproduces the full tree-level string S -matrix (in the massive as well as in the massless sectors). Consider a candidate covariant action for the fields (A, B) . It should of course have the same kinetic terms as in (1). Correspondence with the string theory S -matrix at the 3-point level fixes the 3-couplings of A and B to be the same as in (1) (up to off-shell ambiguities). A covariant action should certainly contain higher order couplings of A 's and A 's and B 's. But (1) is already sufficient for reproducing the correct string S -matrix. Thus all higher order couplings must conspire to give a vanishing contribution to the S -matrix. That this is impossible is clear already from consideration of the 4-point massless amplitude: it is impossible to cancel the contribution of the 4-vertex coming from the

Einstein term (which must be present in a covariant action!) by the contributions of tree diagrams with higher-order couplings ($A^2 B^2$, AB^3 , ...).

The statement of non-existence of a generally covariant action is true even if we "integrate over" all the massive fields except a finite subset B_1 . The reason is that to get an action for A's and B_1 's we are to subtract from the string S-matrix the contributions of exchanges of A's and B_1 's, but a covariant action $S_{\text{cov}}[A]$ is obtained if only the massless exchanges are subtracted out. For example, let us try to construct a covariant action for the metric and the tachyon φ which reproduces all string amplitudes involving gravitons and tachyons. Computing the 3-point amplitudes we find that in addition to the Einstein " h^3 ", φ^3 and $\partial_\mu \varphi \partial_\nu \varphi h_{\mu\nu}$ - couplings there should also be the $\varphi h_{\mu\nu} h_{\mu\nu}$ -coupling. Thus a hypothetical covariant action should look like

$$S[G, \varphi] = c_0 \int d^D x \sqrt{G} \left\{ R + \partial_\mu \varphi \partial_\mu \varphi + m^2 \varphi^2 + c_1 \varphi + c_2 \varphi^3 + \dots \right\} \quad (5)$$

Computing the 4-point graviton amplitude corresponding to (5) we find that in addition to the contribution of the Einstein term there is also the $h^2 (\alpha' m^2)^{-1} h^2$ - contribution of the tachyon exchange. However, the string theory 4-graviton amplitude (the term in it which is of relevant power in momenta) is known (see, e.g. [10]) to be exactly equal to the Einstein 4-point amplitude. Thus the action (5) fails to reproduce the corresponding sector of the string theory S-matrix. It is only when all the massive fields (including the tachyon) are integrated out, the covariant effective action $S_{\text{cov}}[A]$ appears to describe the massless sector of the string S-matrix.

It should be mentioned that it is possible to construct a covariant "action" for all the A- and B-fields if we drop the condition that it should reproduce the string S-matrix in the massive sectors. To find such an "action" we may simply take the string S-matrix with only the massless exchanges subtracted out and replace the on-shell values of the fields by arbitrary off-shell values.

The result will look like $\sum_{cov}$ plus all possible interactions of B's with B's and A's. This "action" is probably equivalent to the "effective action" of ref. [11] (see below). Whether such a covariant "action" for the infinite set of local fields which does not reproduce the string S-matrix is a useful object remains an open question (cf. [12]).

2. Before a discussion of string theory equations of motion within the "6'-model approach" let us demonstrate (on a heuristic level) how the path integral ansatz for the action [11] can be derived from a string field theory (see also [7]). The starting point is the purely "cubic" (unexpanded) form of the string field theory action [5,6] which we write in the following way

$$S = g^{-2} \int \partial c_1 \partial c_2 \partial c_3 \Phi[c_1] \Phi[c_2] \Phi[c_3] K[c_1 c_2 c_3] \quad (6)$$

$$K[c_1 c_2 c_3] = \lim_{\epsilon \rightarrow 0} \int_{M_\epsilon} d\lambda \exp(-I_\epsilon) \quad , \quad I_\epsilon = \frac{1}{4\pi\alpha'} \int d^2\sqrt{g} \partial_\alpha x^\mu \partial^\alpha x^\mu \quad (7)$$

Here $C = \{x^\mu(\epsilon), e(\epsilon)\}$ denotes a closed loop with a one-dimensional metric on it [3] (equivalently, e may be considered as a bosonized ghost [13]). The 3-vertex K is given by a Polyakov path integral over infinitesimal 3-tube surfaces M_ϵ ($\partial M_\epsilon = C_1 \cup C_2 \cup C_3$) [3]. Note that we do not integrate over the two-metric but simply fix a particular metric on M_ϵ as a Weyl gauge (we consider $D = 26$). To rewrite (6) in terms of local fields we shall choose the following "non-standard" parametrization of Φ [7]

$$\Phi[C] = \int_M d\lambda \exp(-I_C) \quad , \quad (8)$$

$$\begin{aligned} I_C = \int d^2\sqrt{g} \Big\{ & \Lambda^2 \varphi(x) + [\partial_\alpha x^\mu \partial^\alpha x^\nu G_{\mu\nu}^{(x)} + R^{(2)} \phi(x) \\ & + i \frac{\epsilon^{\alpha\beta}}{\sqrt{g}} \partial_\alpha x^\mu \partial_\beta x^\nu A_{\mu\nu}^{(x)}] + \Lambda^2 [\partial_\alpha x^\mu \partial^\alpha x^\nu \partial_\beta x^\lambda \partial^\beta x^\rho B_{\mu\nu\lambda\rho}^{(x)} \\ & + \partial^2 x^\mu \partial_\alpha x^\nu \partial^\alpha x^\lambda B_{\mu\nu\lambda}^{(x)} + \partial_\alpha \partial_\beta x^\mu \partial^\alpha x^\nu \partial^\beta x^\lambda \tilde{B}_{\mu\nu\lambda}^{(x)} \\ & + \partial^2 x^\mu \partial^\alpha x^\nu \partial_\alpha x^\lambda B_{\mu\nu}^{(x)} + R^{(2)} \partial_\alpha x^\mu \partial^\alpha x^\nu C_{\mu\nu}^{(x)} + R^{(2)} \partial^2 x^\mu C_\mu^{(x)} + \end{aligned}$$

$$+ (R^{(2)})^2 C(x) + \text{terms with } e^{a\phi}] + O(\Lambda^{-4}) \} \quad (9)$$

Here the integration goes over disc-like two-surfaces ($\partial M = C$). The $\mathcal{G}$ -model action $I_{\mathcal{G}}$ is a most general local reparametrization invariant action of a 2d theory constructed out of x^{μ} and $g_{a\bar{b}}$. It should be stressed that (8) represents an arbitrary functional of C : the set of fields $A' = \{G_{\mu\nu}, \phi, A_{\mu\nu}\}$, $R' = \{\varphi, B_{\mu\nu\lambda\rho}, \dots\}$ can be put in one-to-one correspondence with the set of components (A, B) of Φ in the standard parametrization based on normal mode expansion of $x^{\mu}(\sigma)$ and $e(\sigma)$. To see this one may choose a conformal gauge $g_{a\bar{b}} = e^{2\phi} \delta_{a\bar{b}}$; then (9) reduces to a sum of products of the fields on vertex operators in a $(D+1)$ -dimensional theory on a flat world sheet. Moreover, (8) simply reduces to the standard parametrization to the linear order in the fields $\varphi, G_{\mu\nu} - \delta_{\mu\nu}, \dots$. The exact transformation $(A', B') \rightarrow (A, B)$ is obviously non-linear and non-local. The representation (8) is natural since it is known that the trivial ($G_{\mu\nu} = \delta_{\mu\nu}$) string-theory vacuum functional does have the path integral representation (see e.g. [14,15]). Also, it is likely that a non-trivial string vacuum corresponds to a ground state of a conformal invariant 2d theory [16] and so is given by (8) with $I_{\mathcal{G}}$ taken at a conformal point.

In (9) Λ is a covariant UV cut-off on a world sheet ($g_{a\bar{b}} \Delta \bar{z}^a \Delta z^{\bar{b}} \sim \Lambda^{-2}$) and the fields (taken to be dimensionless as x^{μ}) are "bare" (Λ -dependent). If we substitute (8) into (6) the vertex K glues the three disc-like surfaces together to form a closed sphere-like surface so that we find (assuming that it is possible to ignore the difference between I_0 and $I_{\mathcal{G}}$ on the infinitesimal 3-tube surface)^{1/}

^{1/} The necessity of making this assumption indicates a difficulty with the full background independence in the formulations of refs. [5,6] or in (6), (7) (I_0 in (7) depends on the flat metric).

$$S = g^{-2} \int_{S^2} d^2x \exp(-I_0) \quad (10)$$

Now the integral goes over two-surfaces with a topology of a sphere. Note that (10) in general depends on a choice of a metric on S^2 , i.e. on a choice of a Weyl gauge made in (8). A particular choice of a Weyl gauge corresponds to a particular off-shell representation of the string field theory in terms of local fields. Since (9) is covariant under $x \rightarrow x'(x)$ and the corresponding tensor transformations of the fields the resulting action (10) should also be manifestly covariant. This does not contradict our previous conclusion because (10) need not reproduce the string-theory S-matrix in all its sectors (the transformation $(A', B') \rightarrow (A, B)$ is non-local and so the "equivalence theorem" need not apply). If we set all the fields in $\int_{S^2}$ equal to their trivial "on-shell" values (plane waves on a flat background) then (10) is simply the generating functional for the string-theory S-matrix (on the flat vacuum) [11]. The reason why the local action (and not a non-local off-shell extension of the S-matrix) is obtained in (10) for the massless fields is that the massless exchanges are subtracted out by a 2d renormalization of the G -model [17, 7]. Since the massive exchanges are not subtracted out we get the covariant "action" which was mentioned above. Computing (10) using weak field expansion in powers of the fields we certainly find all the standard 3-point couplings (off-shell extensions of the corresponding amplitudes) being present. In particular, the "action" contains $A'^2 B'$ -couplings (massive "tadpoles") and thus is different from (2) in its dependence on the massive fields. The "kinetic" $B' \Delta(A) B'$ -term in the "action" (representing all string diagrams with two B' -legs, arbitrary number of A' -legs and only massive exchanges) is also different from the $\tilde{B}$ -kinetic term in (2) (representing only massive exchange diagrams without AB^2 -vertices).

3. In view of the vanishing expectation values $\langle V \rangle$ of the vertex operators in a conformal invariant theory on a sphere [18, 19] the classical equations $\frac{\delta S}{\delta g} = 0$, $\frac{\delta S}{\delta G_{\mu\nu}} = 0$, etc., are satisfied if the theory (9) is conformal invariant. Hence the conditions of conformal invariance of the G -model

should be related to string theory equations of motion [17-20]. Within the loop expansion these conditions (vanishing of "β-functions") are manifestly covariant. Loop expansion corresponds to low momentum expansion and is inappropriate for derivation of an action involving massive fields. As was suggested in [21,22] (see also [23]) it is necessary to expand in powers of the massive fields but to sum over all loops. It is clear that this procedure should reproduce the non-covariant action (1) at the quadratic and cubic orders. What happens at higher orders seems to depend on a prescription of how to compute the "β-functions". If we consider high energies (what is necessary if we account for asymptotic massive fields) loop expansion is inappropriate already in the massless sector of the $\mathcal{G}$ -model. A proper treatment of non-perturbative infinities (breaking down the naive covariance of the $\mathcal{G}$ -model) may correspond to subtraction of massive exchanges and may result in equations of motion directly equivalent to that following from (1). I. instead we do loop expansion in the massless sector but sum all the loops with a given number of "massive" insertions the resulting "β-functions" are manifestly covariant and probably equivalent to equations of motion derived from the "action" (10) in which all massive exchanges are not subtracted out. In any case, any prescription that gives covariant equations of motion involving the massive fields as well as the massless ones cannot be (as we discussed above) in correspondence with the string-theory S-matrix.^{2/}

Loop expansion is in general sufficient in order to derive the linear in fields terms in "β-functions". The bare action (9) has the following structure:
$$I_0 = \sum_{n=1}^{\infty} \Lambda^{-2n} \int d^4x \mathcal{O}_n \Phi_n^{(0)},$$
 where $\Phi_n^{(0)}$ are the bare dimensionless fields ($\Phi_1^{(0)} = \varphi$, etc.) and $\mathcal{O}_n$ are polynomials in $\partial_{\mu}\partial_{\nu}x$ and $\partial_{\mu}\partial_{\nu}g$ of dimension $2n+2$ (equal to the number of $\partial_{\mu}\partial_{\nu}x$ in $\mathcal{O}_n$). I_0 is invariant under

^{2/} In particular, the covariant graviton-tachyon action derived in [21] does not contain the $\varphi h_{\mu\nu} h_{\mu\nu}$ -term and thus produces the S-matrix which differs from the string-theory S-matrix already at the 3-point level. We have also checked that the covariantization ($\sqrt{G} \varphi^3$) of the tachyon 3-vertex is not predicted by the string amplitudes with massless and tachyonic exchanges subtracted out (the "subtracted" 5-point amplitude does not contain the $\varphi^3 h_{\mu\nu} h_{\mu\nu}$ -term).

$g_{ab} \rightarrow \lambda^2 g_{ab}$, $z \rightarrow z$, $\Lambda \rightarrow \lambda^2 \Lambda$ ($\lambda = \text{const}$). This invariance is present also in the correlation function; in particular, $\langle \chi(z) \chi(z) \rangle = \frac{1}{4\pi} \ln\left(\frac{\Lambda^2}{\mu^2} e^{2\varphi}\right)$

if $g_{ab} = e^{2\varphi} \delta_{ab}$. Renormalized Φ 's and β -functions are defined as follows

$$\Lambda^{-2n} \Phi_n^{(0)} = \mu^{-2n} \left[\Phi_n + \ln\left(\frac{\Lambda}{\mu}\right) T_n^{(1)}(\Phi) + \dots \right] \quad (11)$$

$$\beta_n = \frac{d\Phi_n}{d\ln\mu} = 2n \Phi_n + T_n^{(1)}(\Phi) \quad (12)$$

The linear in Φ_n term in β_n is easily obtained in one-loop approximation:

$$\langle \Phi_n(y+z) \rangle \rightarrow \frac{i}{2} \partial_\mu \partial_\nu \tilde{\chi}_n \langle \xi^\mu \bar{\xi}^\nu \rangle + \dots, \quad T_n^{(1)} = -\frac{\alpha'}{2} \square \Phi_n + \dots, \\ \beta_n = \frac{\alpha'}{2} (-\square + m_n^2) \Phi_n + \dots, \quad m_n^2 = 4n/\alpha', \quad n = -1, 0, 1, \dots \quad (13)$$

(we have made the rescaling: $G_{\mu\nu} \rightarrow G_{\mu\nu}/4\pi\alpha'$). The local piece of the expectation value of the trace of the energy-momentum tensor T_a^a is given by (cf. [19, 24, 25])

$$\langle T_a^a \rangle_{\text{loc}} = \sum_{n=-1}^{\infty} \mu^{-2n} \mathcal{O}_n(y) \tilde{\beta}_n(y) \quad (14)$$

The $\tilde{\beta}$ -functions are related to β 's and have the following general structure

$$\tilde{\beta}_n = \sum_{m=-1}^{\infty} (-\gamma_{nm} + 2n \delta_{nm}) \tilde{\Phi}_m + O(\Phi^2) \quad (15)$$

where γ is a kind of "anomalous dimension operator" [24, 25]. We expect that $\tilde{\beta}_n = 0$ correspond to $\frac{\delta S}{\delta \Phi_n} = 0$ and hence γ should describe the mixing between the physical and auxiliary fields in the kinetic terms in the action. It was noted in [3] (and also in [24]) that the fields in (9) multiplying the $\mathcal{O}_n$'s which contain the scalar curvature represent the Stueckelberg fields (necessary in order to have local Lorentz-covariant description of higher spin massive fields).

It is possible to check for the first positive (mass)² level that the fields $B_{\mu\nu\lambda\rho}$, $B_{\mu\nu\lambda}$, $\tilde{B}_{\mu\nu\lambda}$, $B_{\mu\nu}$, $C_{\mu\nu}$, C_μ , C in (9) "mix" in the corresponding $\tilde{\beta}_n$ -functions just in the right way. For example, let us show that

$C_{\mu\nu}$ contributes to the $\tilde{\beta}$ -function for $B_{\mu\nu\lambda\rho}$. Varying the $C_{\mu\nu}$ -term in the classical action (9) with respect to the conformal factor of the metric we get: $\mu^{-2} \square_\alpha (\partial_\alpha \chi^\mu \partial^\alpha \chi^\nu C_{\mu\nu})$, or in $\langle T_a^a \rangle$: $\mu^{-2} \partial_\alpha \chi^\mu \partial^\alpha \chi^\nu \partial_\alpha \chi^\lambda \partial^\alpha \chi^\rho \partial_\lambda \partial_\rho C_{\mu\nu}$. Hence $(\Delta \tilde{\beta})_{B_{\mu\nu\lambda\rho}} \sim \partial_\lambda \partial_\rho C_{\mu\nu} + \partial_\mu \partial_\nu C_{\lambda\rho}$. So the action should contain the term

$\int d^2x B^{\mu\nu\lambda\rho} \partial_\mu \partial_\nu C_{\lambda\rho}$. For consistency we are then to find the $\partial_\lambda \partial_\rho B_{\lambda\rho\mu\nu}$ -term in the $\tilde{E}$ -function for $C_{\mu\nu}$. Expanding (9) near $\gamma^{\mu\nu}$

$$T_0 = \int d^2x \sqrt{g} \left\{ \frac{1}{4\pi\alpha'} \partial_\alpha \{ \gamma^{\mu\nu} \} \partial^\alpha \{ \gamma^{\lambda\rho} \} + \dots + \frac{1}{2} \Lambda^{-2} \partial_\alpha \gamma^{\mu\nu} \partial^\alpha \gamma^{\lambda\rho} \partial_\beta \{ \gamma^{\mu\nu} \} \partial^\beta \{ \gamma^{\lambda\rho} \} \partial_\alpha \partial_\beta B_{\mu\nu\lambda\rho}^{(0)} + \dots \right\}$$

we get the two-loop contribution $\sim \mu^{-2} \int d^2x R^{(0)} \alpha_2^{-1} R^{(2)} \partial_\alpha \gamma^{\mu\nu} \partial^\alpha \gamma^{\lambda\rho} \partial_\beta B_{\lambda\rho\mu\nu}$

to the 2d effective action. The corresponding term in $\langle T_a^a \rangle$ is

$$\mu^{-2} R^{(2)} \partial_\alpha \gamma^{\mu\nu} \partial^\alpha \gamma^{\lambda\rho} \partial_\beta B_{\lambda\rho\mu\nu}, \text{ i.e. } (\Delta \tilde{\beta})_{C_{\mu\nu}} \sim \partial_\lambda \partial_\rho B_{\lambda\rho\mu\nu}.$$

In order to correctly establish the

$O(\Phi^2)$ -terms in (15) we are to account for non-perturbative (in α') contributions to the conformal anomaly [21, 22].

In conclusion, we would like to note that the impossibility to describe the dynamics of massive modes in string field theory in a generally covariant way indicates that the ordinary general covariance is only an approximate low-energy symmetry. Thus if compactification in string theory occurs at a string scale ($\sim \sqrt{\alpha'}$) it may be inappropriate to attempt to describe it using the standard concepts of the Einstein's gravitational theory such as curved space-time and generally covariant actions for fields propagating in it.

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