Rolling Tachyon on non-BPS brane.

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Abstract

We consider rolling tachyon on non-BPS brane. Main attention is devoted to a study of equation of the motion obtained within level truncation scheme for the modified fermionic cubic SFT action. Few comments about rolling tachyon matter are also presented.

Contents

1 Introduction

The physical problem which we would like to address in these lectures is a time dependence of the tension of an unstable non-BPS D-brane in the process of its decay. It is expected that starting from its maximal value, equal to the density of the vacuum energy, the tension of an unstable non-BPS D-brane becomes zero when the tachyon field approaches the true vacuum and the D-brane disappears [1]. We start with study of classical time dependent solutions describing a motion of the tachyon on an unstable non-BPS D-brane. One can study the tachyon dynamics within effective Dirac-Born-Infield (DBI) theory [2, 3] closed related with boundary CFT [5, 6, 7] or within open (super)string field theory[8, 9, 10]. In these lectures we discuss mainly the second possibility[11, 12, 13, 14, 16, 17, 18].

Equations of motion describing the tachyons in string field theory have special features. Namely, these equations

- involve infinite number of other fields
- contain infinite number of time derivatives, and hence the initial value problem seems to be not well-defined.

The first obstacle can be avoided

- working in the level truncation scheme proposed by Kostelecky, Samuel [19];
- integrating out all high levels field, that, apparently, is equivalent to working with DBI effective action[20, 21, 22].

To overcome the second obstacle we can just change the problem. One can write equations in the integral form and for these equations consider the boundary problems specifying positions of the tachyon at infinite times. There are some similarities between these equations and equations describing tachyon dynamics for p-adic strings [23]

The tachyon potential and various classical solutions in SFT on an unstable D-brane system have been intensively studied in the last two years (see [1], reviews [8, 9, 10] and references therein). First only time independent solutions representing either the tachyon vacuum or static D-branes of lower dimension have been examined. The interest to time-depending solutions to SFT is related with the desire to study the classical decay process of an unstable D-brane[2]. The process of production and decay of unstable branes can described as a space-like D-brane [24] and also can be studied within a worldsheet analysis [2]. There are also some similarity between tachyon condensation and decay of unstable NC solitons [25].

Here we are going to consider the process of tachyon condensation as a special solution of SFT [11]. There is no problem to construct perturbative solutions of SFT [26], but like usual field theories perturbative solutions cannot describe transitions from one vacuum to other. We will consider the low level truncated NS SFT and show that there is a perturbative solution that interpolates between two non trivial vacua of the theory. At zero time this solution coincides with perturbative vacuum, and therefore can be considered as a solution describing the tachyon condensation.

2 Tachyon Field on D-brane, SFT approach

2.1 Tachyon action and equation of motion in level-truncated SFT

2.1.1 Bosonic string

SFT action is (see notations in reviews [8, 9])

$$S = \frac{1}{2}\langle\langle A, Q_B A\rangle\rangle + \frac{1}{3}\langle\langle A, A, A\rangle\rangle. \tag{2.1}$$

Substituting in (2.1) a string field A only with the tachyon mode one gets [19]

$$S[\phi] = \int \frac{d^{p+1}x}{g_o^2 {\alpha'}^{\frac{(p+1)}{2}}} \left(-\frac{1}{2}\alpha' \partial_\mu \phi(x) \partial^\mu \phi(x) + \frac{1}{2}\phi^2(x) - \frac{1}{3\gamma^3}\tilde{\phi}^3(x)\right), \tag{2.2}$$

$\gamma = \frac{4}{3\sqrt{3}}$. This action describes the selfinteracting tachyon field with cubic potential and infinite number of derivatives in the interaction, since

$$\tilde{\phi}(x) = D_\Box \phi(x) \tag{2.3}$$

where

$$D_\Box = \exp\left(-\alpha' \ln\gamma\, \Box\right)\ , \qquad \Box = \partial_\mu \partial^\mu \tag{2.4}$$

Note that $\ln\gamma < 0$. We use the signature $(-+++\dots)$.

Let us write equation of motion for the action (2.2). We divide the action in to two parts

$$S = \frac{1}{g_o^2 {\alpha'}^{\frac{(p+1)}{2}}}(S_1 + S_2),$$

where S_1 corresponds to the first two summands in (2.2), and S_2 to the third. We want to compute $\delta S/\delta\phi(x)$. We have

$$\frac{\delta S_1}{\delta\phi(x)} = (1+\alpha'\Box)\phi(x)$$

Now let us compute $\delta S_2/\delta\phi(x)$. We have consider an infinitesimal deviation with small filed $h(x)$

$$\delta S_2 = S_2[\phi+h] - S_2[\phi] =$$

$$-\frac{1}{3\gamma^3}\int d^{p+1}x\,\left[(D_\Box\phi - D_\Box h)^3 - (D_\Box\phi)^3\right] \approx$$

$$-\frac{1}{3\gamma^3}\int d^{p+1}x\ 3\ (D_\Box\phi)^2\ D_\Box h(x) =$$

$$-\frac{1}{\gamma^3}\int d^{p+1}x\ (D_\Box(D_\Box\phi)^2)\ h(x)$$

moving to the last equality we used the following property of operator $D_\Box$ which follows from its definition (2.4)

$$\int dx\ D_\Box\varphi(x)\ \phi(x) = \int dx\ \varphi(x)\ D_\Box\phi(x)$$

We have

$$\frac{\delta S_2}{\delta\phi(x)} = -\frac{1}{\gamma^3}\ D_\Box(D_\Box\phi)^2$$

The equations of motion are obtained by equating $\delta S/\delta\phi(x)$ to zero, which gives us

$$(1+\alpha'\Box)\phi(x) - \frac{1}{\gamma^3}\ D_\Box(D_\Box\phi)^2 = 0$$

rewriting it in the form

$$D_\Box^{-1}(1+\alpha'\Box)\phi(x) = \frac{1}{\gamma^3}\ (D_\Box\phi)^2(x)$$

and using (2.3) we finally obtain

$$(1+\alpha'\Box)C_a\tilde{\phi}(x) = \frac{1}{\gamma^3}\ \tilde{\phi}^2(x) \tag{2.5}$$

$$C_a \equiv 2\alpha'\ln\gamma\,\Box, \qquad a \equiv -2\alpha'\ln\gamma$$

On space homogeneous configurations equation of motion is

$$(-\alpha' \frac{d^2}{dt^2} + 1)C_a\tilde{\phi}(t) = \frac{1}{\gamma^3}\ \tilde{\phi}^2(t) \tag{2.6}$$

After rescaling $\Phi = \frac{1}{\gamma^3}\ \tilde{\phi}$ E.O.M. takes the form

$$(-\alpha' \frac{d^2}{dt^2} + 1)C_a\Phi(t) = \Phi^2(t) \tag{2.7}$$

Changing of variables $t \to \alpha t$ gives

$$C_a \to C_{a/\alpha^2}, \qquad \alpha' \frac{d^2}{dt^2} \to \frac{\alpha'}{\alpha^2}\frac{d^2}{dt^2}$$

and (2.7) takes the form

$$(-\underbrace{\frac{\alpha'}{\alpha^2}}_{denote\ q_B^2}\frac{d^2}{dt^2} + 1)C_{a/\alpha^2}\Phi(\alpha t) \quad = \quad \Phi^2(\alpha t) \tag{2.8}$$

Since we want to get the equation of motion with $C \equiv C_{1/4}$ we take α as

$$a/\alpha^2 \equiv 1/4$$

therefore

$$\alpha^2 = 4a, \quad q_B^2 = \frac{\alpha'}{4a} = -\frac{1}{8\ln\gamma}$$

and (2.8) takes the form

$$(-q_B^2\frac{d^2}{dt^2} + 1)C\Phi(\alpha t) = \Phi^2(\alpha t) \tag{2.9}$$

Introducing a new function

$$\Upsilon(t) = \Phi(\alpha t), \tag{2.10}$$

we get the following equation

$$(-q_B^2\frac{d^2}{dt^2} + 1)C\Upsilon(t) = \Upsilon^2(t) \tag{2.11}$$

So, on space homogeneous configurations the equation of motion is

$$(-q_B^2\frac{d^2}{dt^2} + 1)C\Upsilon(t) = \Upsilon^2(t), \quad q_B^2 = -\frac{1}{8\ln\gamma}, \quad C = e^{\frac{1}{4}\partial^2}, \tag{2.12}$$

where $\partial = \partial_t$ and Υ is $\tilde{\phi}$ after rescaling

$$\Upsilon(t) = \frac{1}{\gamma^3}\tilde{\phi}(\alpha t) \tag{2.13}$$

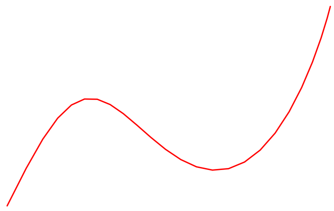


Figure 1: Cubic-potential

2.1.2 Anharmonic oscillator

If one ignores the presence of operator C in the LHS of (2.11) we left with equation of motion

$$(-q^2\partial^2 + 1)x(t) = x^2(t), \tag{2.14}$$

that describes the motion of a particle in the cubic potential.

The cubic potential has two extremum Fig.1. If the initial conditions are such that energy $E = \frac{1}{2}\dot{x}^2 - \frac{1}{2}x^2 + \frac{1}{3}x^3$ is less then zero Fig.2 a) the motion is periodic, if energy is more then zero the motion is unbounded Fig.2 b).

2.1.3 Tachyon action and equation of motion in p-adic string

For $q = 0$ one gets p-adic string equation [23, 12]

$$e^{\frac{1}{4}\partial^2}\Phi = \Phi^2, \tag{2.15}$$

for $p = 2$.

The cubic potential has two extremum and it has been shown [12] that there is no monotonic solution of equation (2.15) interpolating between false and true vacuum. Moreover, in [17] it has been proved that there is no continuous solutions and it has been pointed to the possibility of existence of discontinuous solutions.

2.2 NS string

2.2.1 Tachyon equation of motion in level-truncated SFT

To describe the open string states living on a single non-BPS D-brane one has to consider GSO$\pm$ states [1]. GSO$-$ states are Grassmann even, while GSO$+$ states are Grassmann odd. The unique (up to rescaling of the fields) a)

b)

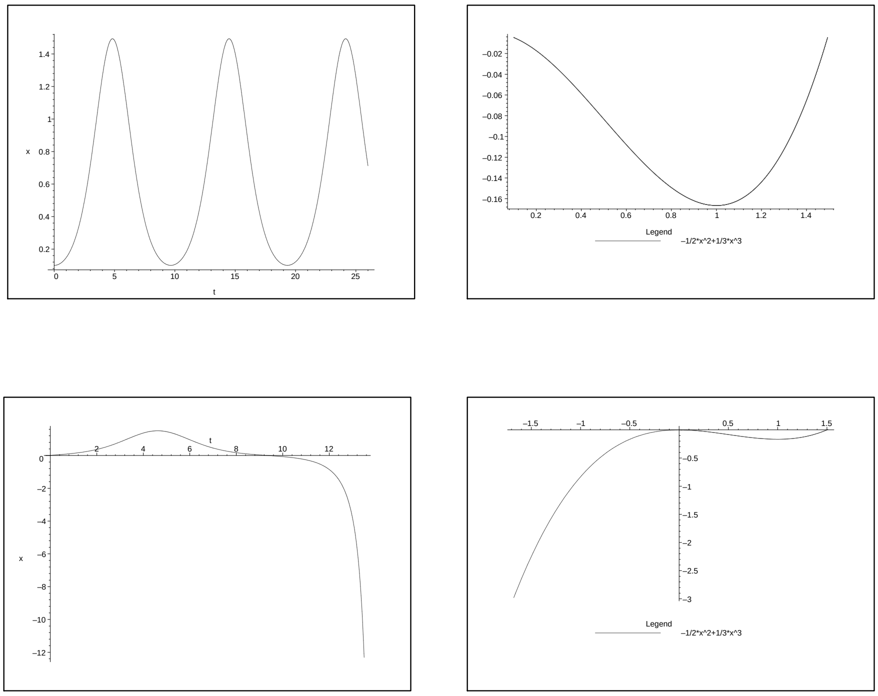


Figure 2: Anharmonic oscillator with cubic-potential; a) periodic trajectory, b) unbounded motion

gauge invariant cubic action unifying GSO+ and GSO− sectors is [27]

$$S = \frac{1}{2}\langle\langle Y_{-2}|A_+, Q_B A_+\rangle\rangle + \frac{1}{3}\langle\langle Y_{-2}|A_+, A_+, A_+\rangle\rangle$$

$$+\frac{1}{2}\langle\langle Y_{-2}|A_-, Q_B A_-\rangle\rangle - \langle\langle Y_{-2}|A_+, A_-, A_-\rangle\rangle.$$

For $A_- = 0$ one gets the action [28, 29]. The low level action contains the tachyon field ϕ and one auxiliary field u

$$S = \int \frac{d^{p+1}x}{g_o^2 \alpha'^{\frac{(p+1)}{2}}} \left[u^2(x) - \frac{\alpha'}{2}\partial_\mu \phi(x)\partial^\mu \phi(x) + \frac{1}{4}\phi^2(x) - \frac{1}{3\gamma^2}\tilde{u}(x)\tilde{\phi}(x)\tilde{\phi}(x) \right], \tag{2.16}$$

where the tilde is as above the action of the operator

$$D_\Box = \exp(-\alpha' \ln \gamma \, \Box), \quad \tilde{\phi} = D_\Box \phi, \quad \text{etc.} \tag{2.17}$$

$\gamma = \frac{4}{3\sqrt{3}}$, $\Box = \partial_\mu \partial^\mu$ and the signature is $(-+++\dots)$. The equations of motion have the form

$$2u(x) - \frac{1}{3\gamma^2} D_\Box \left((D_\Box \phi)^2\right)(x) = 0, \tag{2.18a}$$

$$\left(\alpha' \Box + 1/2\right)\phi(x) - \frac{2}{3\gamma^2} D_\Box \left((D_\Box \phi)(D_\Box u)\right)(x) = 0. \tag{2.18b}$$

Applying $D_\Box$ to both equations above we can write them as

$$\tilde{u} - \frac{1}{6\gamma^2} D_\Box^2 \left((\tilde{\phi})^2\right) = 0, \tag{2.19a}$$

$$\left(\alpha' \Box + 1/2\right)\tilde{\phi} - \frac{2}{3\gamma^2} D_\Box^2 \left(\tilde{u}\tilde{\phi}\right) = 0. \tag{2.19b}$$

Here

$$D_\Box^2 \phi \equiv \exp(-2\alpha' \ln \gamma \, \Box)\phi$$

and for simplicity we do not write arguments of functions. Performing a rescaling of fields and coordinates

$$\Psi(x) = \frac{\sqrt{2}}{3(\gamma)^2}\tilde{\phi}(\alpha x) \quad \text{and} \quad \Upsilon(x) = \frac{4}{3\gamma^2}\tilde{u}(\alpha x) \tag{2.20}$$

with

$$\alpha^2 = 4a = -8\alpha' \ln \gamma \quad \text{and} \quad q^2 = \frac{\alpha'}{2a} = -\frac{1}{4\ln\gamma}. \tag{2.21}$$

we recast the latter system to the following form

$$\Upsilon(x) - D_\Box^2 \left(\Psi^2\right)(x) = 0, \tag{2.22a}$$

$$\left(q^2 \Box + 1\right)\Psi(x) - D_\Box^2 \left(\Upsilon\Psi\right)(x) = 0. \tag{2.22b}$$

Hereafter $D_\Box = \exp(\frac{1}{8}\Box)$.

The equations of motion on space homogeneous configurations are

$$\Upsilon(t) - D^2 \left(\Psi^2\right)(t) \;=\; 0, \tag{2.23}$$

$$\left(-q^2 \partial^2 + 1\right)\Psi(t) - D^2 \left(\Upsilon\Psi\right)(t) \;=\; 0. \tag{2.24}$$

To get these eqs. we make rescaling of fields and time

$$\psi(t) = \frac{\sqrt{2}}{3(\gamma)^2}\phi(\alpha t),\ \Psi(t) = \frac{\sqrt{2}}{3(\gamma)^2}\tilde{\phi}(\alpha t),\ \upsilon(x) = \frac{4}{3\gamma^2}u(\alpha t),\ \Upsilon(t) = \frac{4}{3\gamma^2}\tilde{u}(\alpha t).$$

Operator $D^2 = e^{-\frac{1}{4}\partial^2}$ is not well-defined. However, we can rewrite the system of equations (2.23), (2.24) using only the operator $D^{-2} = e^{\frac{1}{4}\partial^2}$, that admits the integral representation,

$$D^{-2}\Upsilon = \Psi^2, \tag{2.25}$$

$$\left(-q^2 \partial^2 + 1\right) D^{-2}\Psi = \Upsilon\Psi. \tag{2.26}$$

2.2.2 Approximation $u = \tilde{u}$

The approximation $u = \tilde{u}$ is valid when even derivatives of the field u vary slowly. Substituting $u = \tilde{u}$ in action (2.35) and integrating over u and making rescaling we get

$$S[\psi] = \int dt \left[\frac{1}{2}\psi^2(t) + \frac{q^2}{2}\partial\psi(t)^2 - \frac{1}{4}\Psi^4(t) \right], \tag{2.27}$$

here we deal only with time depending configurations and rescaled fields. Equation of motion for this action is

$$\left(-q^2\partial^2 + 1\right) D^{-2}\Psi = \Psi^3. \tag{2.28}$$

The action (2.27) contains infinite number of derivative. As in the case of p-adic string [12] one can rewrite this equations in the integral form

$$\mathcal{C}_{,\ q}\Psi = \Psi^3, \tag{2.29}$$

here and below

$$\mathcal{C}_{a,\ q}\psi(t) = \frac{-q^2\partial^2 + 1}{\sqrt{4\pi a}} \int_{-\infty}^{\infty} e^{-\frac{(t-t')^2}{4a}}\psi(t')dt',$$

where

$$\mathcal{C}_{,\ q} = \mathcal{C}_{1/4,\ q}, \quad \mathcal{C}_a = \mathcal{C}_{a,\ 0}, \quad \mathcal{C} = \mathcal{C}_{1/4}$$

The field ψ is related with Ψ by the integral operator $\psi = \mathcal{C}_{1/8}\Psi$.

2.2.3 Anharmonic oscillator

If we ignore the operator D^{-2} in (2.28) we get equation of motion

$$\left(-q^2\partial^2 + 1\right) \Psi = \Psi^3. \tag{2.30}$$

that for all $q^2 \neq 0$ defines equation of motion for anharmonic oscillator with W-potential. For W-potential all motions are periodic and depending on the total energy are located around one of two minimum (negative total energy), see Fig.3a, or on both minima (positive total energy), see Fig.3b

If we take into account derivatives hidden in D^{-2} and take approximation $D^{-2} \approx 1 + 1/4\partial^2$ instead of eq. (2.28) we get

$$\left((\frac{1}{4} - q^2)\partial^2 + 1 \right) \Psi(t) = \Psi^3(t). \tag{2.31}$$

a)

b)

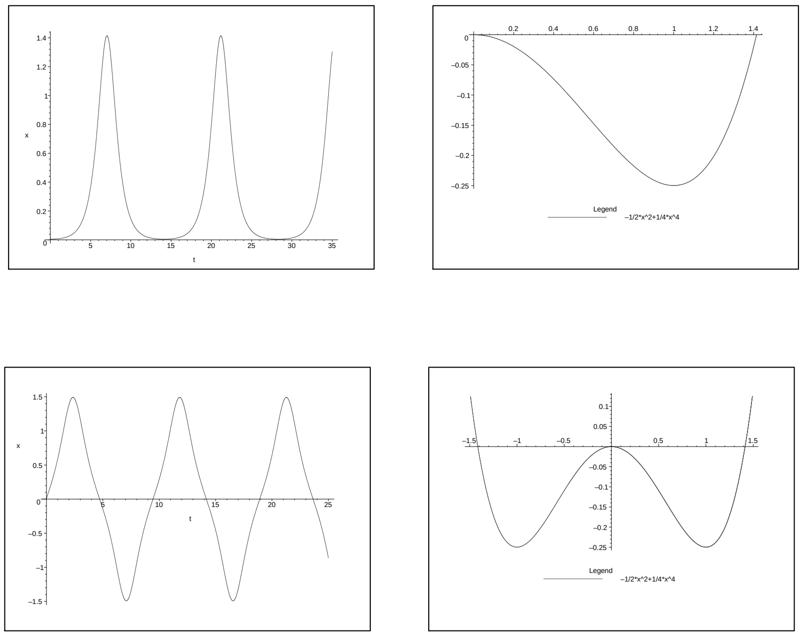


Figure 3: Anharmonic oscillator with W-potential; a) periodic trajectory located around the left minimum, b) periodic trajectory crossing two minima

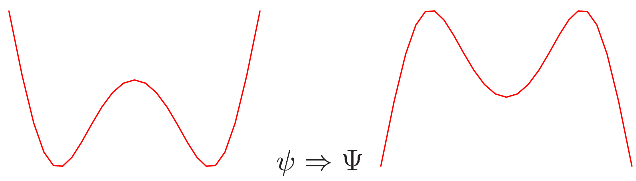


Figure 4: W-potential; -W-potential.

In particular, for $q = 0$ we get an equation describing a particle moving in **-**W-potential Fig.4. Therefore we see that by including of smoothed fields in the action $\psi^4 \to \Psi^4$ we reverse potential for q small enough. For large q we left with W-potential.

This simple observation leads to important physical consequences. In W-potential all motions are periodic, but in **-**W-potential there is a rolling solution, i.e. a solution interpolating between two vacua during infinite large time, i.e. the time requested to rich a vacuum starting from other one is infinite. This rolling solution is just the Wick-rotated version of the well known kink solution.

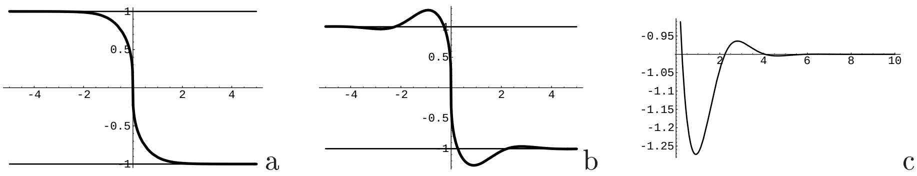


Figure 5: a) $\Psi_0(t)$; b) $\Psi(t)$ for $q^2 = 0.96$; c) oscillations of $\Psi(t)$ with decreasing amplitudes around Ψ_0

2.2.4 Interpolating solution for p-adic string

Equation (2.28) for $q = 0$ is nothing but the p-adic string equation of motion for $p = 3$. This equation has an interpolating solution [23, 12, 14]. We denote this solution by Ψ_0. A rigorous proof of existence of interpolating solution for p-adic equations for odd p is given in [17]. As compare to the kink solution a presence of higher derivatives only changes a profile of Ψ_0, see Fig.5a.

2.2.5 Interpolating solution to eq.(2.28)

To describe interpolating solutions to eq.(2.28) for small q one can consider small fluctuations around Ψ_0, i.e. $\Psi(t) = \Psi_0(t) + \varphi(t)$. One gets equation

$$\left(-q^2\partial^2 + 1\right)\mathcal{C}\varphi = 3\Psi_0^2\varphi + q^2\partial^2\mathcal{C}\Psi_0 + ... \tag{2.32}$$

For $t \to \pm\infty$ $\Psi_0^2 \to 1$ and $\partial^2\mathcal{C}\Psi_0 \to 0$ and we left with a homogeneous integral equation

$$\left(-q^2\partial^2 + 1\right)\mathcal{C}\varphi = 3\varphi \tag{2.33}$$

This equation has solutions $\varphi(t) = Ae^{i\omega t}$ with ω a subject of the following equation

$$\left(q^2\omega^2 + 1\right)e^{-\frac{1}{4}\omega^2} = 3 \tag{2.34}$$

This equation has a complex solution for $q^2 < 1.77$, that corresponds oscillations around Ψ_0 with decreasing amplitude and a real solution for large q^2 that corresponds an appearance of oscillator regime [14].

2.3 Energy and Pressure. SFT approach

To get an expression for the energy-momentum tensor let us write the approximated tachyon action in the form

$$S[\psi] = \int dx\sqrt{-g}\left[\frac{1}{2}\psi^2(x) - \frac{q^2}{2}\partial_\mu\psi(x)\partial^\mu\psi(x) - \frac{1}{4}\Psi^4(x)\right], \tag{2.35}$$

here we include the metric and

$$\Psi(x)=\exp(m\Box)\psi(x),\quad m\equiv\frac{1}{8}\quad\text{and}\quad\Box=\frac{1}{\sqrt{-g}}\partial_{\mu_1}\sqrt{-g}g^{\mu_1\nu_1}\partial_{\nu_1}\quad(2.36)$$

By definition

$$T_{\alpha\beta}(x)=\frac{2}{\sqrt{-g(x)}}\frac{\delta S}{\delta g^{\alpha\beta}(x)}.\qquad(2.37)$$

To find RHS of (2.37) let us find $\frac{\delta S}{\delta g^{\alpha\beta}(x)}$. One has [12]

$$\frac{\delta S}{\delta g^{\alpha\beta}(x)}=\int dy\frac{gg_{\alpha\beta}}{2\sqrt{-g}}\delta(x-y)\left[\frac{1}{2}\psi^2(y)-\frac{q^2}{2}\partial_\mu\psi(y)\partial^\mu\psi(y)-\frac{1}{4}\Psi^4(y)\right]+$$

$$+\int dy\sqrt{-g}(-\frac{q^2}{2}\delta^{\mu\nu}_{\alpha\beta}\partial_\mu\psi(y)\partial_\nu\psi(y)\delta(x-y))-$$

$$-\int dy\sqrt{-g}(\phi+\psi)^3(y)\frac{\delta\phi}{\delta g^{\alpha\beta}}\delta(x-y),$$

where

$$\phi(x)\equiv\Psi(x)-\psi(x)=\sum_{n=0}^{\infty}\frac{m^n}{n!}\Box^n\psi(x)-\psi(x)=\sum_{n=1}^{\infty}\frac{m^n}{n!}\Box^n\psi(x)$$

Therefore

$$\frac{\delta S}{\delta g^{\alpha\beta}(x)}=-\frac{1}{2}\sqrt{-g}\Big[g_{\alpha\beta}(\frac{1}{2}\psi^2(x)-\frac{q^2}{2}\partial_\mu\psi(x)\partial^\mu\psi(x)-\frac{1}{4}\Psi^4(x))+$$

$$+q^2\partial_\alpha\psi(x)\partial_\beta\psi(x)\Big]-\sqrt{-g}\{\frac{g_{\alpha\beta}}{2}\Big[m\int_0^1 d\rho(e^{m\rho\Box}\Psi^3)(\Box e^{m(1-\rho)\Box}\psi)(x)+$$

$$+m\int_0^1 d\rho(\partial_\mu e^{m\rho\Box}\Psi^3)(\partial^\mu e^{m(1-\rho)\Box}\psi)(x)\Big]-$$

$$-m\int_0^1 d\rho(\partial_\alpha e^{m\rho\Box}\Psi^3)(\partial_\beta e^{m(1-\rho)\Box}\psi)(x)\}$$

Note that here we have used the following formula

$$\int dyf(y)\frac{\delta\phi(y)}{\delta g^{\alpha\beta}}=\frac{g_{\alpha\beta}}{2}\Big[m\int_0^1 d\rho(e^{m\rho\Box}f)(\Box e^{m(1-\rho)\Box}\psi)+\qquad(2.38)$$

$$m\int_0^1 d\rho(\partial_\mu e^{m\rho\Box}f)(\partial^\mu e^{m(1-\rho)\Box}\psi)\Big]-m\int_0^1 d\rho(\partial_\alpha e^{m\rho\Box}f)(\partial_\beta e^{m(1-\rho)\Box}\psi)$$

Therefore,

$$T_{\alpha\beta}=-g_{\alpha\beta}(\frac{1}{2}\psi^2(x)-\frac{q^2}{2}\partial_\mu\psi(x)\partial^\mu\psi(x)-\frac{1}{4}\Psi^4(x))-q^2\partial_\alpha\psi(x)\partial_\beta\psi(x)- \tag{2.39}$$

$$-g_{\alpha\beta}\Big[m\int_0^1 d\rho(e^{m\rho\Box}\Psi^3)(\Box e^{m(1-\rho)\Box}\psi)(x)+$$

$$+m\int_0^1 d\rho(\partial_\mu e^{m\rho\Box}\Psi^3)(\partial^\mu e^{m(1-\rho)\Box}\psi)(x)\Big]+$$

$$+2m\int_0^1 d\rho(\partial_\alpha e^{m\rho\Box}\Psi^3)(\partial_\beta e^{m(1-\rho)\Box}\psi)(x)$$

For $\psi=\psi(t)$ we get

$$E(t)=-\frac{1}{2}\psi^2+\frac{q^2}{2}(\partial\psi)^2+\frac{1}{4}\Psi^4- \tag{2.40}$$

$$+m(\int_0^1 d\rho(e^{-m\rho\partial^2}\Psi^3)(\partial^2 e^{-m(1-\rho)\partial^2}\psi)-\int_0^1 d\rho(\partial e^{-m\rho\partial^2}\Psi^3)(\partial e^{-m(1-\rho)\partial^2}\psi)),$$

or

$$E(t)=-\frac{1}{2}\psi^2+\frac{q^2}{2}(\partial\psi)^2+\frac{1}{4}\Psi^4+m\int_0^1 d\rho(e^{-m\rho\partial^2}\Psi^3)\overleftrightarrow{\partial}(e^{-m(1-\rho)\partial^2}\partial\psi) \tag{2.41}$$

One can immediately see that this energy conserves. Indeed,

$$\frac{dE(t)}{dt}=-\psi\partial\psi+q^2(\partial\psi)(\partial^2\psi)+\Psi^3\partial\Psi+m\int_0^1 d\rho(e^{-m\rho\partial^2}\Psi^3)\overleftrightarrow{\partial^2}(e^{-m(1-\rho)\partial^2}\partial\psi)$$

Using the identity

$$-m\int_0^1 d\rho(e^{-m\rho\partial^2}\varphi)\overleftrightarrow{\partial^2}(e^{-m(1-\rho)\partial^2}\phi)=\varphi\overleftrightarrow{e^{-m\partial^2}}\phi, \tag{2.42}$$

we get

$$\frac{dE(t)}{dt}=-\psi\partial\psi+q^2(\partial\psi)(\partial^2\psi)+\Psi^3\partial\Psi+\Psi^3\overleftrightarrow{e^{-m\partial^2}}\partial\psi=$$

$$=-\psi\partial\psi+q^2(\partial\psi)(\partial^2\psi)+\Psi^3\partial\Psi-\Psi^3 e^{-m\partial^2}\partial\psi+\partial\psi e^{-m\partial^2}\Psi^3=$$

$$=-\psi\partial\psi+q^2(\partial\psi)(\partial^2\psi)+\Psi^3\partial\Psi-\Psi^3\partial\Psi+\partial\psi e^{-m\partial^2}\Psi^3=$$

$$= -\partial\psi\{(-q^2\partial^2+1)\psi + e^{-m\partial^2}\Psi^3\} = 0$$

E.O.M. for the action (2.35) is

$$\Psi^3 = e^{1/4\partial^2}(-q^2\partial^2+1)\Psi$$

and energy is

$$E = \mathcal{E}_p + \mathcal{E}_k + \mathcal{E}_{nl}(t), \tag{2.43}$$

i.e. it is the sum of three terms representing kinetic energy

$$\mathcal{E}_k = \frac{q^2}{2}(\partial\psi)^2,$$

potential energy

$$\mathcal{E}_p = -\frac{1}{2}\psi^2 + \frac{1}{4}\Psi^4$$

and non-local terms [13]

$$\mathcal{E}_{nl}(t) = \mathcal{E}_{nl1}(t) + \mathcal{E}_{nl2}(t) = \frac{1}{8}\int_0^1 d\rho \left(e^{\frac{2-\rho}{8}\partial^2}(-q^2\partial^2+1)\Psi\right) \overleftrightarrow{\partial} \left(\partial e^{\frac{\rho}{8}\partial^2}\Psi\right). \tag{2.44}$$

For $q = 0$ both nonlocal terms, $\mathcal{E}_{nl1}(t)$ and $\mathcal{E}_{nl2}(t)$ are negative. Indeed, $\mathcal{E}_{nl1}(t) \approx \frac{1}{8}\Psi_0\partial^2\Psi_0 < 0$, since for negative times $\Psi_0 > 0$ and $\partial^2\Psi_0 < 0$ and for positive times $\Psi_0 < 0$ and $\partial^2\Psi_0 > 0$. $\mathcal{E}_{nl2}(t) \approx -\frac{1}{8}\partial\Psi_0\partial\Psi_0 < 0$.

In Fig.6 we plot functions $\mathcal{E}_p(t) + \mathcal{E}_k(t)$ and $\mathcal{E}_{nl}(t)$. We see that up to inaccuracy in numerical calculations the energy is conserved. To calculate the nonlinear terms numerically we use the integral form

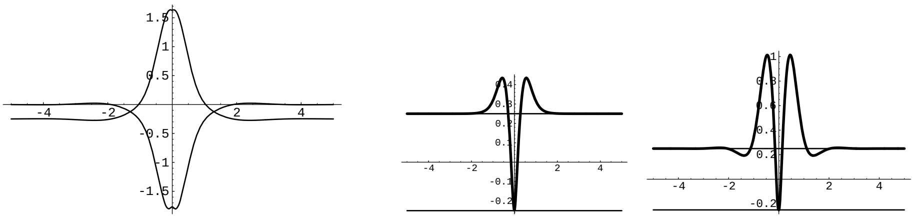
$$\mathcal{E}_{nl}(t) = \frac{1}{8}\int_\varepsilon^1 d\rho(K_1\Psi)(K_2\Psi) - \frac{1}{8}\int_\varepsilon^1 d\rho(K_3\Psi)(K_4\Psi), \tag{2.45}$$

where $K_1 = (-q^2\partial^2+1)C_{\frac{2-\rho}{8}},\ K_2 = \partial^2 C_{\frac{\rho}{8}},\ K_3 = (-q^2\partial^3+\partial)C_{\frac{2-\rho}{8}},\ K_4 = \partial C_{\frac{\rho}{8}}$. In actual calculations we take $\varepsilon \simeq 0.02$.

The pressure is defined in terms of the energy-momentum tensor $P(t)_i = -T_i^i$ (no summation). Since we consider homogeneous configurations we will omit the vector index i. The pressure has the form [13]

$$\mathcal{P}(t) = -E + 2\mathcal{E}_k(t) + 2\mathcal{E}_{nl2}(t), \tag{2.46}$$

The analogous formula takes place for bosonic SFT [12]. For a large time the pressure is equal to $-E$. One can see that the pressure at the zero time is equal to the total energy E, i.e. $\mathcal{P}(0) = E$. This can be shown



a) b) $q^2 = 0$ $q^2 = 0$

Figure 6: a) Testing the constancy of the energy for the interpolating solution to equation (2.29). This figure shows $\mathcal{E}_p(t) + \mathcal{E}_k(t)$ and $\mathcal{E}_{nl}(t)$. Their sum is equal to zero up to numerical errors. b) Calculating the pressure on the solution Fig.5 to equation (2.29) for $q^2 = 0$ and $q^2 = 0.96$. This figure shows $\mathcal{P}(t)$ has minimum at $t = 0$ and goes to $1/4$ when $\pm\infty$

by taking into account the explicit relation between the energy and the pressure (2.46) and the fact that for our solution $\mathcal{E}_{nl1}(0) = 0$.

On Fig.6 b) we show the pressure $\mathcal{P}(t)$ for the solution Ψ_0 as a function of time. We see that in the regions where $\Psi\partial^2\Psi$ is large negative the pressure is large positive, since

$$\mathcal{P}(t) = -E + 2\mathcal{E}_k(t) + 2\mathcal{E}_{nl2}(t) = 2E - 2\mathcal{E}_p - 2\mathcal{E}_{nl1}(t) > 0,$$

that explains two picks in Fig.6b. Two local minima in the figure for $q^2 = 0.96$ correspond oscillations around Ψ_0. The minimum of the pressure is at t=0. For large times the value of $\mathcal{P}$ is positive and equal to $-E = 1/4$.

In the above consideration we take into account only the energy of tachyon field. Following the Sen conjecture one has to add to the action an extra term representing the D-brane tension. This term shifts on a constant the tachyon potential making it equal to zero at its minimum. This constant also shift the pressure and we get that for our solutions the pressure goes to zero for large times.

3 Effective Field Theory Description of Rolling Tachyon

3.1 Boundary states

Initially the decay of D-brave has been statued [2, 3, 3] within a boundary CFT [5, 6, 7]. Boundary CFT is defined as

$$< \mathcal{V}_1 \cdots \mathcal{V}_n >_{D_B} = \int [dx^\mu] \ \ \mathcal{V}_1 \cdots \mathcal{V}_n \ \ e^{-I_{bulk} - I_{bndy}} \tag{3.47}$$

where

$$I_{bulk} = \frac{1}{2\pi}\int_D d^2z g_{\mu\nu}\partial X^\mu \bar{\partial} X^\nu \tag{3.48}$$

D is the unit disk and a boundary interaction is taken in the form

$$I_{bndy} = \lambda \int_{\partial D} dt \, \cosh X^0(t) \tag{3.49}$$

Note that the boundary term

$$I_{bndy} = \lambda \int_{\partial D} dt \, \cos X^0(t) \tag{3.50}$$

defines a solvable boundary conformal field theory [5, 6, 7]. The boundary term (3.49) is the Wick rotated version of (3.50).

A boundary state $|\mathcal{B}>$ associated with D-brane is a closed string state (of ghost # 3) defined so that

$$<\mathcal{B},\mathcal{V}> = <\mathcal{V}>_{D_B}, \tag{3.51}$$

where $|\mathcal{V}>$ is an arbitrary closed state state (of ghost # 3) and $<\mathcal{V}>_{D_B}$ is the one point function of the closed string vertex operator on the unit disk defined by (3.47).

The boundary state for the D25-brane with the boundary interaction (3.49) takes the form

$$|B\rangle = \mathcal{N}_{25}|B\rangle_{X^0} \otimes |B\rangle_{\vec{X}} \otimes |B\rangle_{bc}, \tag{3.52}$$

where $\mathcal{N}_{25}$ is a normalization constant, $|B\rangle_{\vec{X}}$ and $|B\rangle_{bc}$ are the usual boundary states for flat D25-branes:

$$|B\rangle_{\vec{X}} \propto \exp\left(-\sum_{n=1}^{\infty}\frac{1}{n}\alpha^i_{-n}\bar{\alpha}^i_{-n}\right)|0\rangle, \qquad (i=1,\dots,25) \tag{3.53}$$

$$|B\rangle_{bc} \propto \exp\left(-\sum_{n=1}^{\infty}(\bar{b}_{-n}c_{-n} + b_{-n}\bar{c}_{-n})\right)(c_0+\bar{c_0})c_1\bar{c}_1|0\rangle. \tag{3.54}$$

$|B\rangle_{X^0}$ is the part of the boundary state that describes the dynamics of the rolling tachyon. It has been found by Sen [2] (the corresponding boundary state in the Wick rotated theory (3.50) has been constructed in [5, 6, 7] to be

$$|B\rangle_{X^0} = \rho(X^0)|0\rangle + \sigma(X^0)\alpha^0_{-1}\tilde{\alpha}^0_{-1}\mu|0\rangle + ... \tag{3.55}$$

The function ρ, σ are given by

$$\rho(t) = \frac{1}{1+\sin\pi\lambda e^t} + \frac{1}{1+\sin\pi\lambda e^{-t}} - 1, \tag{3.56a}$$

$$\sigma(t) = \cos 2\pi\lambda + 1 - \rho(t) \tag{3.56b}$$

3.2 Energy-momentum tensor

To get the energy-momentum tensor

$$T_{\mu\nu} = -\frac{2}{\sqrt{-g}}\frac{\delta S[g]}{\delta g^{\mu\nu}} \tag{3.57}$$

one can take into account that according $\sigma-$model approach to string theory, the space-time action is given by partition function of world-sheet theory with space-time fields specifying interaction

$$S[g] \propto \langle 1 \rangle_{D_B} \tag{3.58}$$

Therefore the coefficient $\rho(X^0)$ and $\sigma(X^0)$ in (3.55) define the stress tensor of the system

$$T_{00} = \frac{1}{2}T_p(\rho(t) + \sigma(t)), \quad T_{ij} = -T_p\rho(t)\ \delta_{ij} \tag{3.59}$$

Here T_p is the tension of a Dp-brane.

From (3.59) one sees that the pressure goes to zero for $t \to \infty$, i.e. the system becomes pressureless.

3.3 The Dirac-Born-Infeld action

The form of the stress tensor (3.59) with $\rho(t)$ and $\sigma(t)$ given by formulas (3.56a) and (3.56b) can be obtained from the Dirac-Born-Infeld action

$$S_{DBI} = -T_p \int d^{p+1}x V(T)\sqrt{-\det(\eta_{\mu\nu} + \partial_\mu T \partial_\nu T)} \tag{3.60}$$

with the potential $V(T)$

$$V(T) = \frac{T_p}{\cosh T/2} \tag{3.61}$$

(see [2, 34, 35]).

Indeed, the equation of motion for the Dirac-Born-Infeld action (3.60) is

$$\left(g^{\mu\nu} - \frac{\partial^\mu T \partial^\nu T}{1 + (\partial T)^2}\right)\partial_\mu \partial_\nu T = -\frac{V'}{V}(1 + (\partial T)^2) \tag{3.62}$$

The characteristic cones of (3.62) are defined by the inverse metric

$$\left(G^{-1}\right)^{\mu\nu} = g^{\mu\nu} - \frac{\partial^\mu T \partial^\nu T}{1 + (\partial T)^2} \tag{3.63}$$

and do not depend on the potential and strictly spearing the tachyon is no any more the tachyon.

The energy momentum tensor $T^{\mu\nu}$ of the tachyon takes the form

$$T^{\mu\nu} = -V\sqrt{1+(\partial T)^2}\left(G^{-1}\right)^{\mu\nu} \tag{3.64}$$

For homogeneous configurations

$$E = \frac{V(T)}{\sqrt{1-\dot{T}^2}}, \qquad P = -V(T)\sqrt{1-\dot{T}^2} \tag{3.65}$$

For the potential (3.78)

$$\dot{T} = \sqrt{1 - \frac{T_p^2}{E^2\cosh^2\frac{T}{2}}}, \tag{3.66}$$

and therefore,

$$\sinh\frac{T}{2} = \begin{cases} a\cosh\frac{t}{2}, & \frac{T_p^2}{E^2} > 1 \\ a\sinh\frac{t}{2}, & \frac{T_p^2}{E^2} < 1 \end{cases} \tag{3.67}$$

$a = \sqrt{|\frac{T_p^2}{E^2} - 1|}$.

From the boundary CFT part equations (3.59) and (3.56a) give

$$-\frac{P}{T_p} = \rho(t) = \frac{1}{1+\sin\pi\lambda e^t} + \frac{1}{1+\sin\pi\lambda e^{-t}} - 1 = \begin{cases} \frac{T_p}{E}\frac{1}{1+a^2\cosh^2\frac{T}{2}}, & \frac{T_p^2}{E^2} > 1 \\ \frac{T_p}{E}\frac{1}{1+a^2\sinh^2\frac{T}{2}}, & \frac{T_p^2}{E^2} < 1 \end{cases}$$

where the constants are related via

$$a^2\sin\pi\lambda = (\frac{T_p}{E} - 1)^2.$$

Using (3.67) one sees that

$$-\frac{P}{T_p} = \frac{T_p}{E}\frac{1}{\cosh^2\frac{T}{2}} \tag{3.68}$$

that is in accordance with

$$P = -\frac{V^2(T)}{E} \tag{3.69}$$

((3.69) follows from (3.65)) for potential given by (3.61).

3.4 Boundary states in superstring theory

The boundary state describing the rolling of the tachyon on unstable D-brane systems in superstring theory can also be obtained in a similar way [2, 3]. One considers a pair of D9 and $\bar{\text{D}}9$ branes in type IIB theory or a single non-BPS D9-brane in type IIA theory. The D9-$\bar{\text{D}}9$ brane pair is described by the boundary CFT with 2×2 Chan-Paton factors and with an appropriate GSO projection [1].

For this case a tachyon field is taken of the form

$$T(x^0) \propto \cosh(x^0/\sqrt{2}), \tag{3.70}$$

Taken into account a Chan-Paton factor σ_1 one gets following the boundary interaction

$$\tilde{\lambda} \int dt\, \psi^0(t) \sinh\left(\frac{X^0(t)}{\sqrt{2}}\right) \otimes \sigma_1, \tag{3.71}$$

and corresponding boundary states

$$|B\rangle = |B,+\rangle - |B,-\rangle, \tag{3.72}$$

where

$$|B,\epsilon\rangle \propto |B,\epsilon\rangle_{X^0,\psi^0} \otimes |B,\epsilon\rangle_{\vec{X},\vec{\psi}} \otimes |B,\epsilon\rangle_{ghost}, \qquad (\epsilon = \pm). \tag{3.73}$$

The spatial and the ghost parts take the usual expressions:

$$\begin{aligned}
|B,\epsilon\rangle_{\vec{X},\vec{\psi}} &\propto \exp\left(-\sum_{n=1}^{\infty} \frac{1}{n} \alpha^j_{-n} \bar{\alpha}^j_{-n}\right) \exp\left(-i\epsilon \sum_{n=1}^{\infty} \psi^j_{-n-1/2} \bar{\psi}^j_{-n-1/2}\right) |0\rangle, \\
|B,\epsilon\rangle_{ghost} &\propto \exp\left(-\sum_{n=1}^{\infty} (\bar{b}_{-n} c_{-n} + b_{-n} \bar{c}_{-n})\right) \\
&\quad \times \exp\left(-i\epsilon \sum_{n=1}^{\infty} (\bar{\beta}_{-n-1/2} \gamma_{-n-1/2} - \beta_{-n-1/2} \bar{\gamma}_{-n-1/2})\right) \\
&\quad \times (c_0 + \bar{c_0}) c_1 \bar{c}_1 e^{-\phi(0)} e^{-\bar{\phi}(0)} |0\rangle.
\end{aligned} \tag{3.74}$$

and X^0, ϕ^0 part has the form

$$|B,\epsilon\rangle_{X^0,\psi^0} = \left(\hat{\rho}(X^0(0)) + i\epsilon \psi^0_{-1/2} \bar{\psi}^0_{-1/2} \hat{\sigma}(X^0(0))\right) |0\rangle, \tag{3.75}$$

where

$$\hat{\rho}(t) = \frac{1}{1 + \sin^2(\pi\tilde{\lambda}) e^{\sqrt{2}t}} + \frac{1}{1 + \sin^2(\pi\tilde{\lambda}) e^{-\sqrt{2}t}} - 1, \tag{3.76}$$

$$\hat{\sigma}(t) = \cos(2\pi\tilde{\lambda}) + 1 - \hat{f}(t). \tag{3.77}$$

The stress tensor of the system is given by (3.59) with $\rho \to \hat{\rho}$ and $\sigma \to \hat{\sigma}$ and $\hat{\rho}$ and $\hat{\sigma}$ given by (3.76) and (3.77). From (3.59) and (3.76) and (3.77) one once again sees that the system becomes pressureless for large t.

The form of the stress tensor (3.59) with $\hat{\rho}(t)$ and $\hat{\sigma}(t)$ given by formulas (3.76) and (3.77) can be obtained from the Dirac-Born-Infeld action (3.60) with potential $V(T)$

$$V_s(T) = \frac{1}{\cosh T/\sqrt{2}} \tag{3.78}$$

3.5 Few comments about cosmological applications

Recently the string tachyon condensation has been discussed in cosmology in context of the observed acceleration of the Universe and the problem of the dark energy. Only 0.3 part of total energy density of the Universe, $\sim 10^{-120}M_p^4 = 10_{-29}g/cm$ is the ordinary matter, the rest part of the energy density corresponds to the dark energy. One of possible scenarios describing the evolution of the Universe filled by the dark energy is based on consideration of gravity interacting with the tachyon matter given by the action (3.60) (see review [36] and references therein).

Since for homogeneous configurations we have (3.65) the state equation reads

$$P = \frac{V^2(T)}{E} \tag{3.79}$$

If for large t $V(T) \to -A = const$ then $P = -\frac{A}{E}$ and one gets the Chaplygin gas.

The Freedman metric is defined as a solution of the gravity equation with the tachyon matter defined by the action (3.60),

$$ds^2 = -dt^2 + a^2(t)d\vec{x}^2 \tag{3.80}$$

and

$$\frac{\ddot{a}}{a} = \frac{\kappa}{3}\frac{V(T)}{\sqrt{1-\dot{T}^2}}(1 - \frac{3}{2}\dot{T}^2) \tag{3.81}$$

The expanding Universe is accelerating at domain $T(t) \approx 0$, $|\dot{T}| \ll 1$ since the RHS of (reff-eq=a) is positive define. As time t goes, T exceeds $\frac{3}{2}$ so that expansion of the Universe slows down.

4 Discussion and Conclusion

In these lectures we have discussed the SSFT tachyon classical time dependent equations of motion describing non-stability of non-BPS D-branes. There is no conventional initial value problem for these equations. We have discussed the special time dependent solution to this equation that describes the rolling tachyon. It starts from the perturbative vacuum and approaches one of stable vacua in infinite time. The pressure on this solution tends to zero for large times. These results have been obtained in the first non-trivial level within the level truncation method and under assumption of a special approximation $u = \tilde{u}$. The first non-trivial level without this assumption has been studied in [13, 14]. It would be very interesting to take into account higher level fields including vector fields and study a dynamics of a D-brane tension.

We have also briefly reviewed the tachyon effective field theory description of non-BPS D-branes and its relation with boundary CFT. A recent study of this subject has suggest that during the decay process the D-brane energy is converted to closed string [37, 16, 38, 18]. It would be interesting to study this problem within SFT for the case of non-BPS brane

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