Bouncing universe with the non-minimally coupled scaler field and its reconstructing

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Abstract

In this paper we consider a non-minimally coupled scaler field, and show its equation of state parameter can crossing over $-1$, $\omega \rightarrow -1$, and bouncing condition. Also we obtain the stability conditions and consider reconstructing for our model.

Keywords: Bouncing Universe; Stability Condition; Coupled Scaler Fields; Reconstructing.

1 Introduction

Scalar fields play a central role in the modern cosmology. The combined analysis of the type Ia supernovae, galaxy clusters measurements and WMAP data provides an evidence for the accelerated cosmic expansion \cite{1-5}. The cosmological acceleration strongly indicates that the present day universe is dominated by smoothly distributed slowly varying Dark Energy (DE) component. The modern constraints on the DE state parameter are around the cosmological constant value, $\omega = -1 \pm 0.1$ \cite{3-7} and a possibility that $\omega$ is varied in time is not excluded. From the theoretical point of view there are three essentially different cases: $\omega > -1$ (quintessence), $\omega = -1$ (cosmological constant) and $\omega < -1$ (phantom) \cite{8-28} and Refs. therein).

Since from the observational point of view there is no barrier between these three possibilities it is worth to consider models where these three cases are realized. Under general

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assumptions it is proved in [29] that within one scalar field model one can realize only one possibility: \( \omega \geq 1 \) (usual model), or \( \omega \leq 1 \) (phantom model). It is interesting that the interaction with the cold dark matter does not change the situation and does not remove the cosmological constant barrier [30].

On the other hand, the Friedman equation forms the starting point for almost all investigations in cosmology. Over the past few years possible corrections to the Friedman equation have been derived or proposed in a number of different contexts, generally inspired by braneworld investigation [31, 32]. These modification are often of a form that involves the total energy density \( \rho \). In [33], multi-scalar coupled to gravity is studied in the context of conventional Friedman cosmology. It is found that the cosmological trajectories can be viewed as geodesic motion in an augmented target space.

In this way, there are several phenomenological models describing the crossing of the cosmological constant barrier [34]-[43]. Most of them use more then one scalar field or use a non-minimal coupling with the gravity, or modified gravity, in particular via the brane-world scenarios. In two-field models one of these two fields is a phantom, other one is a usual field and the interaction is nonpolynomial in general. It is important to find a model which follows from the fundamental principles and describes a crossing of the \( \omega = -1 \) barrier.

In this paper we show that such a model may appear within a brane approach when the universe is considered as a slowly decaying D3-brane and a possibility to cross the barrier comes from taking into account a back reaction of the D3-brane. This DE model [24] assumes that our universe is a slowly decaying D3-brane and its dynamics is described by the open string tachyon mode and the back reaction of this brane is incorporated in the dynamics of the closed string tachyon. The open string tachyon dynamics is described within a level truncated open string field theory (OSFT). The notable feature of this OSFT description of the tachyon dynamics is a non-local polynomial interaction [34]-[47].

It turns out the open string tachyon behavior is effectively described by a scalar field with a negative kinetic term (phantom) [48]-[51]. However this model does not suffer from quantum instability, which usually phantom models have, since in the nonlocal theory obtained from OSFT there are no ghosts at all near the non-perturbative vacuum [24].

On the other hand, inflation [52]-[54] is possibly the only known mechanism which dynamically solves the flatness and the horizon problem of the universe. Thus it has become an almost indispensable ingredient in cosmology. The inflaton, a scalar field, can also produce the density perturbations causally which can match with the data from observation. For example, the recent WMAP data [55]-[58] strongly supports the idea that the early universe went through an inflationary phase. Usually one considers the inflationary phase to be driven by the potential of a scalar field. Recently there has been an upsurge in activity for constructing such models in string theory. In the context of string theory, the tachyon field in the world volume theory of the open string stretched between a D-brane and an anti-D-brane or on a non-BPS D-brane has been taken as a natural candidate to play the role of the inflaton [59, 60]. This possibility of the tachyon field driving the cosmological inflation is related to the decay of unstable brane as a time dependent process which was advocated by Sen [61, 62]. The effective action used in the study of tachyon cosmology consists of the standard Einstein-Hilbert action and an effective action for the tachyon field on unstable
D-brane or braneantibrane system.

In this way, there are a lot of cosmological observations, such as SNe Ia [63]-[66], WMAP [67, 68], SDSS [69, 70, 71], Chandra X-ray Observatory [72] etc., that they reveal some cross-checked information of our universe. They suggest that the universe is spatially flat, and consists of approximately 70% dark energy with negative pressure, 30% dust matter (cold dark matters plus baryons), and negligible radiation, and that the universe is undergoing an accelerated expansion.

To accelerate the expansion, the equation of state parameter \( \omega \equiv \frac{p}{\rho} \) of the dark energy must satisfy \( \omega < -\frac{1}{3} \), where \( p \) and \( \rho \) are its pressure and energy density, respectively. The simplest candidate of the dark energy is a tiny positive time-independent cosmological constant \( \Lambda \), for which \( \omega = -1 \). Another possibility is quintessence [73]-[77], a cosmic real scalar field that is displaced from the minimum of its potential.

Briefly, in section 2 we have defined our extended model and we have shown the equation of state parameter must be crossing over \( -1, \omega \to -1 \), and the Hubble parameter \( H \) running across zero at \( t = 0 \) and ingratiates bouncing conditions. In section 3 we consider stability of our model. In section 4 we will reconstruct a non-minimally scaler filed in the three forms of parametrization. Finally in section 5 we compare them together, with numerical methods.

2 Non-minimally coupled scalar field

In this section we consider the action in the Jordan frame [78, 79] with \( g_{\mu\nu} \) metric:

\[
S = \int d^4 x \sqrt{-g} \left[ \frac{M_p^2}{2} R - f(\phi) R - \frac{1}{2} (\nabla \phi)^2 - V(\phi) \right],
\]

where \( M_p^2 \equiv \frac{1}{8\pi G} \), and \( f(\phi)R \) term corresponds to the non-minimal coupling of the scalar field to gravity. For a flat Friedman-Robertson-Walker (FRW) universe we can assume that the universe is described by the flat, homogeneous, and isotropic universe model with the scale factor \( a \). With (1) we obtain the equation of motion of the scalar field \( \phi \),

\[
\ddot{\phi} + 3H \dot{\phi} + V' + 6f'(H + 2H^2) = 0.
\]

Where \( V' = \frac{dV(\phi)}{d\phi} \) and \( f' = \frac{df(\phi)}{d\phi} \), and \( H = \frac{\dot{a}}{a} \) is the Hubble parameter. The energy-momentum tensor \( T^{\mu\nu} \) is given by the standard:

\[
\delta g_{\mu \nu} S = - \int d^4 x \frac{\sqrt{-g}}{2} T^{\mu \nu} \delta g_{\mu \nu}.
\]

We can read the energy density from (3) as

\[
\rho = \frac{1}{2} \dot{\phi}^2 + V + 6H(f + Hf),
\]
and similarly the energy pressure

\[ p = \frac{1}{2} \dot{\phi}^2 - V - 2\ddot{f} - 4H\dot{f} - 2f(2\dot{H} + 3H^2). \]  

(5)

With the Friedman equations,

\[ H^2 = \frac{1}{3M_p^2} \rho, \]  

(6)

and

\[ \dot{H} = -\frac{1}{2M_p^2}(\rho + p), \]  

(7)

The continuity equation to be,

\[ \dot{\rho} + 3H \rho (1 + \omega) = 0, \]  

(8)

and from Eqs. (11) and (10) we have

\[ H = \frac{1}{6} \sqrt{12 \left( \dot{f}(3\dot{f} - \dot{\phi}^2) + V(M_p^2 - 2f) \right) + 6M_p^2 \dot{\phi}^2}. \]  

(9)

We now study the cosmological evolution of equation of state for the present model. The equation of state is \( p = \omega \rho \). To explore the possibility of the \( \omega \) across \(-1\), we have to check \( \frac{d}{dt}(\rho + p) \neq 0 \) when \( \omega \to -1 \). From Eqs. (11) and (5) one can obtain the following expressions,

\[ \rho + p = \dot{\phi}^2 + 2H\dot{f} - 4\dot{H}f - 2\ddot{f}. \]  

(10)

Therefore from Eqs. (7) and (10) we have

\[ \dot{H} = -\frac{1}{2(M_p^2 - 2f)}(\dot{\phi}^2 + 2H\dot{f} - 2\ddot{f}), \]  

(11)

so we have \( \frac{d}{dt}(\rho + p) = -2M_p^2 \dot{H} \) or,

\[ \dot{H} = \frac{1}{(M_p^2 - 2f)} \left( \frac{d\dot{f}}{dt} + \dot{H}\dot{f} - H\ddot{f} - \dot{\phi}\ddot{\phi} \right). \]  

(12)

Now with replacing Eqs. (9) and (11) in the Eq. (12), if \( f = 1 + \sum_{i=1} c_i \phi^{2i} \) and a tachyon scaler field, \( V = V_0 e^{-\lambda \phi^2} \) (motivated by string theory [80]), with \( c_1 = \frac{1}{12} \) and \( c_i = 0 \) for \( i > 1 \),
one can obtain

\[
\ddot{H} = \frac{1}{(6M_p^2 - \phi^2)} \left\{ \frac{d\dot{\phi}}{dt} [\dot{\phi} - \frac{6M_p^2\dot{\phi}}{U}] + \frac{\ddot{\phi}[3\dot{\phi} + \frac{6M_p^2(\lambda\phi V - 2) + 2\phi V(1 - \lambda\phi^2)}{U} + \frac{6M_p^2\phi V(\lambda(\phi^3 - 6M_p^2) - \phi)}{U^3}}}{6M_p^2 - \phi^2} \left( \frac{(2\phi U - 3M_p^2\phi)\phi^2}{U} + (\phi - U)\phi \right) \right\] 
+ \frac{\ddot{\phi}[2V]}{U} \left( 2\lambda(3M_p^2 - 4\lambda\phi^2)M_p^2 + \phi(1 - \lambda\phi^2(5 - \lambda\phi^2)) \right) + \frac{\phi}{(6M_p^2 - \phi^2)} \times \left( \frac{\phi^2}{U}(4U + 2V(1 + \lambda(6M_p^2 - \phi^2))) + 2\phi U(\frac{4\phi^2}{(6M_p^2 - \phi^2)} + 1) \right) \right\] 
+ \frac{V^2(\lambda(\phi^3 - 6M_p^2) - \phi)^2 + 9M_p^4\phi^2\phi^2}{U}. \tag{13}
\]

Where \( U = \sqrt{2(6M_p^2 - \phi^2)V + 6M_p^2\phi} \). With implying Eqs. (10) and (11) and a lot of complexity calculations one can finds either (i) \( \dot{\phi} = 0 \) or (ii) \( \ddot{\phi} = 0 \) when \( \omega \rightarrow -1 \).

(i) Let us assume \( \dot{\phi} = 0 \) first when \( \omega \rightarrow -1 \), so we have,

\[
\ddot{H} = \frac{1}{(6M_p^2 - \phi^2)} \left\{ \frac{d\dot{\phi}}{dt}[\dot{\phi} - \frac{6M_p^2\dot{\phi}}{U}] + \sqrt{2(6M_p^2 - \phi^2)V} \left( 6M_p^2(\lambda\phi V - 2) - 2\phi V(1 + \lambda\phi^2) \right) \right\} 
+ \frac{V^2(\lambda(\phi^3 - 6M_p^2) - \phi)^2}{U}. \tag{14}
\]

Therefore the conditions for having the \( \omega \) across over \( -1 \) are (i-1) \( \frac{d\dot{\phi}}{dt} \neq 0 \), (i-2) \( \ddot{\phi} \neq 0 \) and (i-3) \( \phi^2 \neq 6M_p^2 \) in addition to the \( \dot{\phi} = 0 \).

(ii) Let us turn to the second case, when \( \omega \rightarrow -1 \). We now have

\[
\ddot{H} = \frac{1}{(6M_p^2 - \phi^2)} \left\{ \frac{d\dot{\phi}}{dt}[\dot{\phi} - \frac{6M_p^2\dot{\phi}}{U}] + \frac{\ddot{\phi}[2V]}{U} \left( 2\lambda(3M_p^2 - 4\lambda\phi^2)M_p^2 + \phi(1 - \lambda\phi^2(5 - \lambda\phi^2)) \right) + \frac{\phi}{(6M_p^2 - \phi^2)} \times \left( \frac{\phi^2}{U}(4U + 2V(1 + \lambda(6M_p^2 - \phi^2))) + 2\phi U(\frac{4\phi^2}{(6M_p^2 - \phi^2)} + 1) \right) \right\] 
+ \frac{V^2(\lambda(\phi^3 - 6M_p^2) - \phi)^2}{U}. \tag{15}
\]

Therefore the conditions for having the \( \omega \) across over \( -1 \) are (ii-1) \( \frac{d\dot{\phi}}{dt} \neq 0 \), (ii-2) \( \ddot{\phi} \neq 0 \) and (ii-3) \( \phi \neq 6M_p^2 \) in addition to the \( \dot{\phi} = 0 \).
In the Fig. 1 we consider a tachyon scaler field, $V(\phi) = V_0 e^{-\lambda \phi^2}$ and $f(\phi) = 1 + \sum_{i=1}^k c_i \phi^{2i}$. We have shown $\omega \to -1$, which gives rise to a possible inflationary phase after the bouncing. Now we consider a detailed examination on the necessary conditions required for a successful bounce. During the contracting phase, the scale factor $a(t)$ is decreasing, i.e., $\ddot{a} < 0$, and in the expanding phase we have $\ddot{a} > 0$. At the bouncing point, $\ddot{a} = 0$, and around this point $\dddot{a} > 0$ for a period of time. Equivalently in the bouncing cosmology the Hubble parameter $H$ runs across zero from $H < 0$ to $H > 0$ and $H = 0$ at the bouncing point. A successful bounce requires around this point,

$$\dot{H} = -\frac{1}{2M_p^2} (\rho + p) = -\frac{1}{2M_p^2} \rho (1 + \omega) > 0. \quad (16)$$

Fig. 2 shows that the Hubble parameter $H$ running across zero at $t = 0$ and ingratiates bouncing conditions, Eq. (16). In the Fig. 2 is seen in which the $H$ crossing zero axes and it is a bounce point.

### Stability conditions

In this section, we deal on the stability of our model. Here we want to consider the stability by a useful the function $c_s^2 = dp/d\rho$. This function just is the stability of our system by the scalar field that it must become more than zero. Of course this function express sound speed in a prefect liquid. Therefore we apply the analysis of our model in the $\omega' - \omega$ plane which $\omega' = d\omega/d \ln a$. In that case we obtain from Eqs. (9) and as the following form,

$$c_s^2 = \frac{\rho'}{\rho} = \omega + \frac{\rho}{\rho'} \omega', \quad (17)$$
Figure 2: The graph of Hubble parameter $H$, for the potential $V(\phi) = V_0 e^{-\lambda \phi^2}$ which $V_0 = 4$, $\lambda = 0.6$, $c_1 = \frac{1}{12}$ and $c_{i>1} = 0$. Initial values are $\phi(0) = 0.5$ and $\dot{\phi}(0) = -1.55$.

where the prime denote derivation with respect $\ln a$. By substituting Eq. (17) in term the prime in above equation yields,

$$c_s^2 = \omega - \frac{\omega'}{3(1 + \omega)}.$$  \hfill (18)

To employing the stability condition ($c_s^2 > 0$), we obtain the two regions $\omega > -1$ and $\omega < -1$ respectively for $\omega' > 3\omega(1 + \omega)$ and $\omega' < 3\omega(1 + \omega)$ in the $\omega' - \omega$ phase plane. These regions have been showed in Fig. 3.

Figure 3: The graph of $\omega'$ with respect to $\omega$ for the two cases investible.

Now we show the stability condition for non-minimally coupled scalar field. We consider the stability condition by $c_s^2$ parameter, which it can rewrite as the Hubbel parameter,

$$c_s^2 = -1 - \frac{\ddot{H}}{3HH}.$$  \hfill (19)
Now by numerical computing we can plot the $c_s^2$, in term of time evolution, that is shown in Fig. 4. We can see the stability condition for late time evolution i.e. $c_s^2$ for $t \rightarrow +\infty$.

Figure 4: The graph of $c_s^2$ with respect to time evolution.

4 Reconstructing a non-minimally coupled scalar field

In this section we consider reconstruct a non-minimally scalar filed in the three forms of parametrization.

Therefore from Eqs. (4) and (5) one can finds,

\[ 2\rho + 3p = \frac{5M_p^2}{2(M_p^2 - 2f)} \dot{\phi}^2 - \frac{M_p^2}{(M_p^2 - 2f)} V - \frac{6M_p^2}{(M_p^2 - 2f)} \ddot{f} = -\bar{K}, \]

(20)

with comparing by Eq. (11) and implying Eq. (9) we have,

\[ \rho = \bar{K} + 3\bar{V}. \]

(21)

Where

\[ \bar{V} = \frac{M_p^2}{(M_p^2 - 2f)^2} \left( 2\dot{f} + (\dot{\phi}^2 - 2\ddot{f})(M_p^2 - 2f) \right) \]

\[ - \frac{M_p^2 \dot{f}}{3(M_p^2 - 2f)^2} \sqrt{36\dot{f}^2 + 6(M_p^2 - 2f)\dot{\phi}^2 + 12V(M_p^2 - 2f)}. \]

(22)

In this method the energy pressure is,

\[ p = -(\bar{K} + 2\bar{V}), \]

(23)

and equation of state can be obtained such as following,

\[ \omega = -1 + \left( \frac{\bar{V}}{\bar{K} + 3\bar{V}} \right). \]

(24)
Then we can rewrite the Friedman equations from Eqs. (6) and (7) as following,

$$3M_p^2 H^2 = \rho_m + \rho = \rho_m + \ddot{K} + 3\dot{V},$$  \hspace{1cm} (25) 

and

$$2M_p^2 \dot{H} = -(\rho_m + \rho + p) = -\rho_m - \ddot{V}. $$ \hspace{1cm} (26) 

Where $\rho_m$ is the energy density of dust matter, therefore we have

$$\ddot{K} = 2\rho_m + 3M_p^2 H^2 + 6M_p^2 \dot{H},$$ \hspace{1cm} (27) 

and

$$\ddot{V} = -\rho_m - 2M_p^2 \dot{H}.$$ \hspace{1cm} (28) 

In this model, the dark energy fluid does not couple to the background fluid, the expression of the energy density of dust matter in respect of redshift $z$ is

$$\rho_m = 3M_p^2 H_0^2 \Omega_{m0} (1 + z)^3,$$ \hspace{1cm} (29) 

where $\Omega_{m0}$ is the ratio density parameter of matter fluid and the subscript 0 indicates the present value of the corresponding quantity. Using the following relation

$$\frac{d}{dt} = -H (1 + z) \frac{d}{dz},$$ \hspace{1cm} (30) 

one can rewrite $\ddot{K}$ and $\ddot{V}$ as following

$$\ddot{K} = M_p^2 H_0^2 \left( 6 \Omega_{m0} (1 + z)^3 + 3r - 3(1 + z)r^{(1)} \right),$$ \hspace{1cm} (31) 

and

$$\ddot{V} = M_p^2 H_0^2 \left( -3 \Omega_{m0} (1 + z)^3 + (1 + z)r^{(1)} \right).$$ \hspace{1cm} (32) 

Where $r = \frac{H^2}{H_0^2}$ and $r^{(n)} = \frac{d^n r}{dz^n}$. By using Eqs. (24), (31) and (32) we obtain following expression for equation of state,

$$\omega = \frac{(1 + z)r^{(1)} - 3r}{-3\Omega_{m0} (1 + z)^3 + 3r},$$ \hspace{1cm} (33) 

and with using Eqs. (21), (23) and (30) one can find the sound speed such as following,

$$c_s^2 = \frac{(1 + z)r^{(2)} - 2r^{(1)}}{-9\Omega_{m0} (1 + z)^2 + 3r^{(1)}},$$ \hspace{1cm} (34) 

the sound speed is discussed for investigation of stability of the model and it necessary is to be $c_s^2 \geq 0$. 

Finally from Eq. (33) we can obtain following equation for $r(z)$

$$r(z) = \Omega_{m0} (1 + z)^3 + (1 - \Omega_{m0}) e^{\int_0^z \frac{1+\omega(\tilde{\tau})}{1+z} \tilde{d} \tilde{\tau}}$$ \hspace{1cm} (35) 

Also by using Eq (32) we have an expression for deceleration parameter $q$ as follows,

$$q = -1 - \frac{\dot{H}}{H^2} = \frac{(1 + z)r^{(1)} - 2r}{2r}. $$ \hspace{1cm} (36)
5 Parametrization

Now act three different forms of parametrization as following and compare them together, with numerical methods.

**Parametrization 1:**
First parametrization has proposed by Chevallier and Polarski [82] and Linder [83], where the EoS of dark energy in term of redshift \( z \) is given by,

\[
\omega(z) = \omega_0 + \frac{\omega_a z}{1 + z}.
\]  

**Parametrization 2:**
Another the EoS in term of redshift \( z \) has proposed by Jassal, Bagla and Padmanabhan [84] as,

\[
\omega(z) = \omega_0 + \frac{\omega_b z}{(1 + z)^2}.
\]  

**Parametrization 3:**
Third parametrization has proposed by Alam, Sahni and Starobinsky [85]. They take expression of \( r \) in term of \( z \) as followoing,

\[
r(z) = \Omega_{m0}(1 + z)^3 + A_0 + A_1(1 + z) + A_2(1 + z)^2.
\]  

By using the results of Refs.[86, 87, 88], we get coefficients of parametrization 1 as \( \Omega_{m0} = 0.29, \omega_0 = -1.07 \) and \( \omega_a = 0.85 \), coefficients of parametrization 2 as \( \Omega_{m0} = 0.28, \omega_0 = -1.37 \) and \( \omega_b = 3.39 \) and coefficients of parametrization 3 as \( \Omega_{m0} = 0.30, A_0 = 1, A_1 = -0.48 \) and \( A_2 = 0.25 \).

The evolution of \( \omega(z) \) and \( q(z) \) are plotted in Fig. 5 and Fig. 6 respectively. Also, using Eqs. (31), (32) and the three parameterizations, the evolutions of \( K(\tilde{z}) \) and \( V(\tilde{z}) \) are shown in Fig. 7 and Fig. 8 respectively.

![Figure 5](image-url)

**Figure 5:** Graphs for the EoS parameter in respect of redshift \( z \). The solid, dot and dash lines represent parametrization 1, 2 and 3 respectively.
The evolution of a scalar field in respect of redshift $z$ are same for all of parametrization 1, 2 and 3. Therefore all of parametrization give us consequence well. Fig. 5 show us in which we can explicitly see the dynamics of a scalar field. Also slope of graph decrease in the early epoch.

Now to achieve to stability of the model by using Eq. (34), so we can obtain following condition for all of parametrization

$$r(z) \geq \Omega_{m0}(1 + z)^3,$$

where is accurate for three above parametrization.

### 6 Conclusion

In this paper, we have investigated bouncing universe by the action in the Jordan frame metric with Eq. (1). In this form has coupled a functional of scalar field with gravity term. We obtained equation of state with respect to time evolution, and one has crossing $\omega \rightarrow -1$. By consider stability, we have drown speed of sound in term of time evolution and it has
Figure 8: Graphs for the reconstructed $\tilde{V}$ in respect of redshift $z$. The solid, dot and dash lines represent parametrization 1, 2 and 3 respectively.

Figure 9: Graphs for the reconstructed $\tilde{V}$ in respect of scaler field $\phi$. The solid, dot and dash lines represent parametrization 1, 2 and 3 respectively.

shown existence of stability in late time. In the graph of Hubble parameter, we can see bouncing condition as $H$ crossing from $t = 0$ i.e. $\dot{H} > 0$.

Also, we have reconstructed all of cosmology parameters with respect to redshift $z$. In that case, we have described the EoS to cross from $-1$ for three parametrization. Therefore, we can see from graph of $z$, the second parametrization is better than two other parametrization.

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