D-brane as Dark Matter in Warped String Compactification

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Abstract

It is pointed out that in the warped string compactification, motion of D-branes near the bottom of a throat behaves like dark matter. Several scenarios for production of the dark matter are suggested, including one based on the D/\bar{D} interaction at the late stage of D/\bar{D} inflation.

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1 Introduction

M/string theory \cite{1,2} is considered as a strong candidate for a unified theory of fundamental physics. Its mathematical consistency and beauty have been attracting interest of many physicists. On the other hand, one of the drawbacks is lack of direct experimental or observational evidence of such a structure at high energies. Having this situation, it seems rather natural to turn our eyes to cosmology and look for cosmological implication of M/string theory since the universe is supposed to have experienced a high energy epoch at its early stage.

Considering the success of inflation \cite{3,4,5} as a scenario of the early universe and the observational evidences for accelerating expansion of the present universe \cite{6,7,8}, one of the important steps toward establishment of M/string cosmology would be construction of de Sitter or quasi-de Sitter universe. However, for a long time it seemed rather difficult to construct 4-dimensional de Sitter universe in M/string theory, especially if we seriously take account of the moduli stabilization. Indeed, the no-go theorem of \cite{9,10} says that in a large class of supergravity theories, there is no no-singular (warped) compactification to 4-dimensional de Sitter space with a finite 4-dimensional Newton’s constant.

There are several proposals to evade the no-go theorem by inclusion of additional sources such as stringy corrections to the supergravity theories in the g_s or \alpha’ expansion and extended sources, i.e. branes. The recent proposal by Kachru, Kallosh, Linde and Trivedi (KKLT) \cite{11} utilizes various ingredients of string theory including warped geometry, fluxes, D-branes, \bar{D}-branes, instanton corrections to moduli potential in order to construct a meta-stable de Sitter vacua in string theory \footnote{See \cite{12} and \cite{13,14} for other proposals of de Sitter and transiently accelerating universe.}.

In the KKLT setup \bar{D}3-branes play an essential role. Inclusion of \bar{D}3-branes at the bottom of a warped throat uplifts stable AdS vacua with negative cosmological constant to meta-stable de Sitter vacua with positive cosmological constant in a theoretically controllable way. Without \bar{D}3-branes, we would end up with a negative cosmological constant, which is inconsistent with observations.
The purpose of this paper is to point out that motion of the $\bar{D}3$-branes near the bottom of a warped throat behaves like dark matter and that it can be naturally generated in the context of brane cosmology. This paper is organized as follows. In Sec. 2 we briefly review the Klebanov-Strassler geometry which approximates the geometry in a warped throat, and consider a probe $\bar{D}3$-brane. In Sec. 3 we promote the $\bar{D}3$-brane action formulated in a 4-dimensional flat background to a curved background and investigate its implications to cosmology. In particular, we show that a $\bar{D}3$-brane near the bottom of a warped throat behaves like dark matter. In Sec. 4 we suggest several possible scenarios to produce the dark matter. Finally, Sec. 5 is devoted to a summary of this paper and discussion.

## 2 $\bar{D}3$-brane in Klebanov-Strassler geometry

In the KKLT setup [11] a throat region is described by the warped deformed conifold solution of Klebanov and Strassler [16]. To be precise, the warped deformed conifold geometry is compactified by additional fluxes as described by Giddings, Kachru and Polchinski [15]. Hence, the geometry in the UV region of the throat significantly deviates from the Klebanov-Strassler solution, but the geometry near the bottom of the throat, i.e. in the IR region, is approximated by the Klebanov-Strassler solution.

The Klebanov-Strassler geometry [16, 17] has the simple ansatz:

$$ds^2 = h^{-1/2}(\tau)\eta_{\mu\nu}dx^\mu dx^\nu + h^{1/2}(\tau)ds_6^2,$$

where $x^\mu (\mu = 0, \cdots, 3)$ are 4-dimensional coordinates and $ds_6^2$ is the metric of the deformed conifold [18, 19]

$$ds_6^2 = \frac{\epsilon^{4/3}K(\tau)}{2} \left[ \frac{1}{3K^3(\tau)} (d\tau^2 + (g^7)^2) + \cosh^2 \left( \frac{T}{2} \right) ((g^3)^2 + (g^4)^2) + \sinh^2 \left( \frac{T}{2} \right) ((g^1)^2 + (g^2)^2) \right].$$

Here,

$$K(\tau) = \frac{(\sinh(2\tau) - 2\tau)^{1/3}}{2^{1/3} \sinh \tau},$$

and $g^i (i = 1, \cdots, 5)$ are orthonormal basis defined by

$$g^1 = \frac{e^1 - e^3}{\sqrt{2}}, \quad g^2 = \frac{e^2 - e^4}{\sqrt{2}},$$

$$g^3 = \frac{e^1 + e^3}{\sqrt{2}}, \quad g^4 = \frac{e^2 + e^4}{\sqrt{2}}, \quad g^5 = e^5,$$

where

$$e^1 \equiv - \sin \theta_1 d\phi_1, \quad e^2 \equiv d\theta_1,$$

$$e^3 \equiv \cos \psi \sin \theta_2 d\phi_2 - \sin \psi d\theta_2,$$

$$e^4 \equiv \sin \psi \sin \theta_2 d\phi_2 + \cos \psi d\theta_2,$$

$$e^5 \equiv d\psi + \cos \theta_1 d\phi_1 + \cos \theta_2 d\phi_2.$$

Because of the warp factor $h^{-1/2}(\tau)$, this geometry is often called the warped deformed conifold. The R-R 3-form field strength $F_3$ and the NS-NS 2-form potential $B_2$ also have the $Z_2$ symmetric $((\theta_1, \phi_1) \leftrightarrow (\theta_2, \phi_2))$ ansatz:

$$F_3 = \frac{M\alpha'}{2} \left\{ g^5 \wedge g^3 \wedge g^4 + d[F(\tau)(g^3 \wedge g^3 + g^2 \wedge g^4)] \right\},$$

$$B_2 = \frac{g_s M\alpha'}{2} \left\{ f(\tau)g^1 \wedge g^2 + k(\tau)g^3 \wedge g^4 \right\},$$

(6)
where $F(0) = 0$ and $F(\infty) = 1/2$. For this ansatz with the additional condition

$$g_s^2 F_3^2 = H_3^2,$$  \hfill (7)

we can consistently set the dilaton $\phi$ and the R-R scalar $C_0$ to zero. The BPS saturated solution found by Klebanov and Strassler \cite{16} is

$$F(\tau) = \frac{\sinh \tau - \tau}{2 \sinh \tau},$$
$$f(\tau) = \frac{\tau \coth \tau - 1}{2 \sinh \tau} (\cosh \tau - 1),$$
$$k(\tau) = \frac{\tau \coth \tau - 1}{2 \sinh \tau} (\cosh \tau + 1),$$ \hfill (8)

and

$$h(\tau) = 2^{2/3} \cdot (g_s M \alpha')^2 e^{-8/3} I(\tau),$$ \hfill (9)

where

$$I(\tau) = \int_{\tau}^{\infty} dx \frac{x \coth x - 1}{\sinh^2 x} (\sinh(2x) - 2x)^{1/3}. \hfill (10)$$

For this solution,

$$C_4 = h^{-1} dx^0 \wedge dx^1 \wedge dx^2 \wedge dx^3$$ \hfill (11)

in a particular gauge. For large $g_s M$ the curvature is small everywhere and we can trust the supergravity description.

When $g_s M$ is sufficiently large, we can treat a $\bar{D}3$-brane as a probe brane \cite{20}. The action for the probe $\bar{D}3$-brane is

$$S_{\bar{D}3} = -T_3 \int d^4 \xi e^{-\phi} \sqrt{-\det(G_{\alpha\beta} - B_{\alpha\beta})} - T_3 \int d^4 \xi C_4,$$ \hfill (12)

where $\xi^\alpha (\alpha = 0, \ldots, 3)$ are intrinsic coordinates on the $\bar{D}3$-brane, $T_3$ is the tension and

$$G_{\alpha\beta} = G_{MN} \frac{\partial x^M}{\partial \xi^\alpha} \frac{\partial x^N}{\partial \xi^\beta}, \quad B_{\alpha\beta} = (B_2)_{MN} \frac{\partial x^M}{\partial \xi^\alpha} \frac{\partial x^N}{\partial \xi^\beta}. \hfill (13)$$

In the following we shall adopt a gauge in which brane coordinates $\xi^\alpha$ coincide with $x^\alpha$:

$$x^\alpha = \xi^\alpha, \quad \psi^m = \psi^m(\xi^\alpha),$$ \hfill (14)

where $\{\psi^m\} (m = 5, \ldots, 10)$ represents $\{\tau, \psi, \theta_1, \phi_1, \theta_2, \phi_2\}$. In the non-relativistic limit,

$$S_{\bar{D}3} = T_3 \int d^4 \xi \left[ -\frac{1}{2} \gamma_{mn} \eta^{\alpha\beta} \frac{\partial \psi^m}{\partial \xi^\alpha} \frac{\partial \psi^n}{\partial \xi^\beta} - 2 h^{-1} \right]. \hfill (15)$$

where $\gamma_{mn} d\psi^m d\psi^n = ds_6^2$ is the metric of the deformed conifold given by \cite{2}.  

### 3 Cosmological implication

We have considered a probe $\bar{D}3$-brane in the Klebanov-Strassler geometry and have obtained the effective action \cite{15}. Let us now consider cosmological implication of this action \cite{2}. With moduli stabilization, the 4-dimensional Einstein gravity should be recovered \cite{3}. In cosmological

\footnote{Another approach to cosmology with the warped string compactification would be to consider time-dependent extension of the Klebanov-Strassler solution. However, without compactification \cite{16} and moduli stabilization \cite{11}, there are instabilities associated with the unfixed moduli \cite{21}.}

\footnote{Evidence can be seen for the recovery of the 4-dimensional Einstein gravity in a simplified setup of warped flux compactification \cite{22}.}
situations, we still expect that the KS geometry is a good approximation to the geometry of extra dimensions in the throat region as far as the energy scale associated with the 4-dimensional physics is sufficiently lower than the stabilization scale. Hence, we promote the action formulated in the flat 4-dimensional spacetime to a curved background $S_{D3} = T_3 \int \sqrt{-g}^{(4)} d^4 \xi \left[ -\frac{1}{2} \gamma_{mn}^{(4)} \frac{\partial \psi_{\alpha}}{\partial \xi^n} \frac{\partial \psi_{\beta}}{\partial \xi^m} - 2h^{-1} \right]$. (16)

We expect that this action describes the dynamics of the brane at energy scales sufficiently lower than the stabilization scale. In general, stabilization of the volume modulus introduces corrections to the effective action, in particular to the mass of the fields. Actually, in Ref. 23 it was pointed out that in the context of $D3/\bar{D}3$ inflation, the inflaton mass receives corrections of order $H$ due to the modulus stabilization. Those corrections were significant for the inflaton dynamics since the inflaton mass must be fine-tuned to a value much smaller than $H$ in order to realize a sufficient inflation. On the other hand, our interest in this paper is on a massive field (which shall be denoted by $\psi$ below). As far as the mass of the field of interest is much larger than the Hubble expansion rate $H$, the corrections due to the modulus stabilization should be ignorable.

With the FRW ansatz

$$ds^2 = -dt^2 + a(t)^2 dx^2,$$ (17)

the equation of motion of $\psi^m$ is

$$\ddot{\psi} + 3H \dot{\psi} + \frac{\partial}{\partial \psi^m} \rho(\psi, \pi) = 0,$$ (18)

where a dot denotes the time derivative,

$$\pi_m = T_3 \gamma_{mn} \ddot{\psi}^n,$$ (19)

and

$$\rho(\psi, \pi) = \frac{1}{2T_3} \gamma_{mn} \pi_m \pi_n + 2T_3 h^{-1}$$ (20)

is the energy density written in terms of $\psi^m$ and $\pi_m$.

Hereafter, we consider the $D3$-brane near the bottom of the throat $\tau = 0$. Thus, we take the small $\tau$ limit of the deformed conifold metric:

$$\gamma_{mn} d\psi^m d\psi^n = \frac{4^{1/3}}{2} \left( \frac{2}{3} \right)^{1/3} \left\{ \frac{1}{2} d\tau^2 + \left[ \frac{1}{2} (g^3)^2 + (g^3)^2 + (g^4)^2 \right] + \frac{1}{4} \tau^2 \left[ (g^1)^2 + (g^2)^2 \right] \right\},$$ (21)

where the first square bracket represents the regular $S^3$ at the bottom of the throat and the second square bracket represents the shrinking $S^3$. With this form of the deformed conifold metric, it is shown by using the equations of motion that there are constants of motion $J_3^1$ and $J_2^1$ related to the angular momenta along the $S^3$ and the $S^2$, respectively. For later convenience the constants of motion $J_3^1$ and $J_2^1$ are normalized as

$$\left( \frac{2}{4^{1/3} T_3} \right)^2 \left( \frac{3}{2} \right)^{1/3} \frac{J_3}{2a^6} = \frac{1}{2} (g^3(\partial_3))^2 + (g^3(\partial_3))^2 + (g^4(\partial_3))^2,$$

$$\left( \frac{8}{4^{1/3} T_3} \right)^2 \left( \frac{3}{2} \right)^{2/3} \frac{J_2}{2a^6} = (g^1(\partial_1))^2 + (g^2(\partial_2))^2.$$ (22)

It is shown that $a^3 \pi_{\phi_1}$ and $a^3 \pi_{\phi_2}$ are also constants of motion, but they do not appear in the following discussions. The equation of motion for $\tau$ and the energy density $\rho$ are written in terms of $J_3^1$ and $J_2^1$ as

$$\ddot{\tau} + 3H \dot{\tau} + \left( \frac{\partial V}{\partial \tau} \right) = 0,$$ (23)

$$\rho = \frac{1}{2} \dot{\tau}^2 + V + \rho_0 + \rho_0,$$ (24)
where
\[
\varphi = \sqrt{\frac{\epsilon^{3/3} T_3}{2}} \left( \frac{2}{3} \right)^{1/6} \tau,
\]
\[
V(\varphi, a) = \frac{1}{2} m^2 \varphi^2 + \frac{J_2^2}{2a^6 \varphi^2},
\]
\[
m^2 = \frac{8}{\epsilon^{3/3}} \left( \frac{3}{2} \right)^{1/3} \frac{d^2}{d\tau^2} h^{-1} \bigg|_{\tau=0} = \frac{4 \cdot 2^{2/3}}{3 H(0)^2} \left( g_s M \alpha' \right)^2.
\] (25)

and
\[
\rho_6 = \frac{1}{\epsilon^{3/3} T_3} \frac{J_2^2}{2a^6},
\]
\[
\rho_0 = 2 T_3 \ h^{-1} \bigg|_{\tau=0}.
\] (26)

Here, \( I(0) \simeq 0.71805 \).

For \( J_2 \neq 0 \), the potential \( V \) is minimized if \( \varphi \) evolves as
\[
\varphi = \sqrt{\frac{|J_2|}{m}} \frac{1}{a^{3/2}}.
\] (27)

Actually, this is a solution to the equation of motion (23) if the FRW universe evolves as \( H \propto a^{-3/2} \). (We shall confirm this cosmological evolution below.) With this solution, the energy density \( \rho \) is
\[
\rho = \left( 1 + \frac{9 H^2}{8 m^2} \right) \rho_3 + \rho_6 + \rho_0,
\] (28)

where
\[
\rho_3 = \frac{|mJ_2|}{a^3}.
\] (29)

At low energy we can safely neglect \( 9 H^2 / 8 m^2 \) compared to 1. The term \( \rho_6 \) can also be neglected since it decays faster than other terms as the universe expands. Therefore, we obtain
\[
\rho \simeq \rho_3 + \rho_0.
\] (30)

The first term behaves like dark matter while the second term contributes to the cosmological constant. In this paper we shall not try to solve the cosmological constant problem and simply assume that \( \rho_0 \) is (almost) canceled by other contributions to the cosmological constant (eg. stabilization of the volume modulus, tension of other branes, etc.). With this assumption, the cosmological energy density is dominated by \( \rho_3 \) and it is confirmed that the cosmological evolution is \( H \propto a^{3/2} \) as assumed, provided that the Friedmann equation is recovered at low energy. Now let us consider a small perturbation around the solution (27) for \( J_2 \neq 0 \):
\[
\varphi = \sqrt{\frac{|J_2|}{m}} \frac{1}{a^{3/2}} + \delta \varphi(t).
\] (31)

The linearized equation of motion implies that the perturbation \( \delta \varphi \) behaves like a massive scalar field with mass \( 2m \),
\[
\delta \ddot{\varphi} + 3H \delta \dot{\varphi} + (2m)^2 \delta \varphi = 0,
\] (32)

and that the solution (27) is indeed stable. At low energy \( H \ll m \), the stress energy tensor due to \( \delta \varphi \) behaves like a pressure-less dust (\( \propto a^{-3} \)) if it is averaged over a timescale sufficiently longer than \( m^{-1} \) and sufficiently shorter than \( H^{-1} \).

For \( J_2 = 0 \), \( \varphi \) behaves as a massive scalar field with mass \( m \) and at low energy \( H \ll m \), its energy density behaves like a pressure-less dust (\( \propto a^{-3} \)) if it is averaged over a timescale sufficiently longer than \( m^{-1} \) and sufficiently shorter than \( H^{-1} \).

In summary we have shown that the stress energy tensor due to the motion of a D3-brane near the bottom of a warped throat region behaves like dark matter \( \rho \propto a^{-3} \). It is of course possible to consider motion of more than one D3-branes as multi-component dark matter.
4  Dark matter production

Now let us discuss possible scenarios to generate the motion of $D$-branes as dark matter.

(i) $D$/$\bar{D}$ interaction at late stage of brane inflation: In Ref. [23] the attractive force between a $D$-brane and a $\bar{D}$-brane were considered as the origin of the inflaton potential. $D$-branes are placed at the bottom of a throat and a mobile $D$-brane falls toward the bottom. This process inevitably generates motion of not only the mobile $D$-brane but also all $\bar{D}$-branes near the bottom of the throat since the $D$/$\bar{D}$ interaction acts between the mobile $D$-brane and each $\bar{D}$-brane. If the number of $\bar{D}$-branes is more than one then there remain remnant $\bar{D}$-branes after the brane inflation followed by the annihilation of one $D$/$\bar{D}$ pair and reheating.

(ii) Closed string modes in the bulk: The motion of $D$-branes may be generated at reheating after inflation if physics in the bulk of extra dimensions plays important roles in the reheating process. For example, in Ref. [24] it was suggested that in the context of the $D$/$\bar{D}$ brane inflation, reheating on our brane placed at the bottom of a throat can be due to closed string modes in the bulk produced during tachyon condensation. In this case the closed string modes in the bulk can kick $D$-branes at the bottom and generate their motion.

(iii) Gravitational production: It may also be possible to generate the dark matter within the context of a 4-dimensional effective theory without thinking about extra dimensions. Gravitational production had been considered as a mechanism to produce superheavy dark matter (often called WIMPZILLA) [25]. The same mechanism should work for the production of the $D$ dark matter if a 4-dimensional effective theory is valid at the time of production.

In all cases, it is expected that the total angular momentum along the $S^2$ of the Klebanov-Strassler geometry should decay to an extremely small value already during inflation, but if there are more than one $D$-branes then each $D$-brane can obtain angular momentum via the above mechanisms. (Of course the total angular momentum remains extremely small.) In this case $D$-branes may collide at late time and lose their angular momenta since the sum of their angular momenta are essentially zero. This may lead to intriguing cosmological and astrophysical consequences such as heavy particle production followed by lepto- and/or baryo-genesis, generation of ultra high energy cosmic rays, impact on the nucleo synthesis, etc. Further studies along this line are worthwhile.

5  Summary and discussion

We have pointed out that in the warped string compactification, motion of $D$-branes near the bottom of a throat behaves like dark matter. The existence of this type of dark matter is a necessary consequence of the KKLT setup, where $D$-branes are required to uplift an AdS spacetime to a de Sitter spacetime.

Our arguments so far have been based on the implicit assumption that our brane is somewhere in the compact manifold where the warp factor is of order unity. This is the reason why we have written down the effective action for the $D$-brane in terms of the 4-dimensional metric $g^{(4)}_{\mu\nu}$. Indeed, the physical (or induced) metric on our brane in this setup is $g_{\mu\nu}$ up to an overall constant factor of order unity. The dark matter mass $m$ is given in [25] and can be estimated as

$$m^2 = \frac{2(4\pi)^{7/4}}{3^{3/2}(0)} \cdot \frac{h^{−1/2}(0)}{g^{1/2} M} \approx 3 \times h^{−1/2}(0) \cdot \left( \frac{g_s}{0.1} \right)^{1/2} \cdot \left( \frac{g_s M}{10} \right)^{−1}, \quad (33)$$

where we have used the formula $M_{10}^2 = 2/(2\pi^7 \alpha' M^2)$ for the 10D gravity scale $M_{10}$. Hence, if we suppose the TeV gravity, $M_{10} \sim TeV$, then the dark matter mass is lighter than TeV by the factor $h^{−1/4}(0)$. Note that the warp factor $h^{−1/4}(0)$ is given by

$$h^{−1/4}(0) \approx \exp \left( \frac{−2\pi K}{3g_s M} \right), \quad (34)$$

and can be exponentially small, where the positive integer $K$ is the value of NS-NS flux required for moduli stabilization [13].
On the other hand, if our brane is located near the bottom of the throat then the dark matter mass should be different. The field $\varphi$ is no longer canonically normalized with respect to the physical (or induced) metric $\bar{g}^{(4)}_{\mu\nu} = h^{-1/2}(0)g^{(4)}_{\mu\nu}$, and the canonically normalized field is $\bar{\varphi} \equiv h^{1/2}(0)\varphi$. Correspondingly, the physical mass of the field $\bar{\varphi}$ is $\bar{m} = h^{1/4}(0)m$ and, thus,
\begin{equation}
\frac{\bar{m}^2}{M^2_{10}} \simeq 3 \times \left(\frac{g_s}{0.1}\right)^{1/2} \cdot \left(\frac{g_s M_{10}}{10}\right)^{-1}.
\end{equation}
Therefore, in this case the dark matter mass is around $\text{TeV}$ if we adopt the TeV gravity.

We have suggested several scenarios for the dark matter production. One of them is based on the $D/\bar{D}$ interaction at the late stage of $D/\bar{D}$ inflation. Suppose that inflation in our 4-dimensional universe is driven by the modulus representing the position of a mobile $D$-brane relative to $\bar{D}$-branes near the bottom of a throat. As argued in ref. [23], a successful inflation in this setup is possible although fine-tuning is required to ensure the flatness of the inflaton potential. The rolling of the inflaton is due to an attractive force between the branes, and the same force inevitably generates the motion of $\bar{D}$-branes as well. Since the $D$-brane potential has the minimum at the bottom of the throat, the $\bar{D}$-branes should start moving around the bottom. When the inter-brane distance becomes sufficiently short, the inflation ends and a pair of $D$- and $\bar{D}$-branes annihilates. The 4-dimensional universe should be reheated by the energy released in the annihilation process. After the annihilation of a $D/\bar{D}$ pair and the reheating of the universe, the remaining $\bar{D}$-branes are still moving around the bottom. As we have shown in this paper, the $\bar{D}$-brane motion around the bottom behaves like dark matter in the 4-dimensional universe. In this context, the flatness of the inflaton potential requires that our brane should be located somewhere in the Calabi-Yau manifold where the warp factor $h^{-1/2}(\tau)$ is of order unity $^4$. This means that the dark matter mass should be much lighter than the 10$D$ gravity scale, say $\text{TeV}$, as discussed above.

In order for this scenario to work, it must be made sure that the amount of produced dark matter is not too much. Otherwise, the energy density of the dark matter ($\propto a^{-3}$) would start dominating the radiation energy density ($\propto a^{-4}$) too soon. In the above example, the production is due to the $D/\bar{D}$ interaction and the interaction is very weak by definition: the enough inflation requires a flat potential and the flat potential implies the weak interaction between branes. From this consideration, it is expected that the amount of dark matter (i.e. the amplitude of $\bar{D}$-brane motion) after inflation should be very small. Very importantly, this expectation seems compatible with the fine-tuning required for successful inflation since both the smallness of the production rate and the flatness of the inflaton potential are related to the weakness of the inter-brane interaction. It is certainly worthwhile quantifying how much fine-tuning this requires.

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