Weak gravity in the Dvali-Gabadadze-Porrati braneworld model

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We analyze weak gravity in the braneworld model proposed by Dvali, Gabadadze, and Porrati, in which the unperturbed background spacetime is given by a five-dimensional Minkowski bulk with a brane which has an induced Einstein-Hilbert term. This model has a critical length scale \( r_c \). Naively, we expect that four-dimensional general relativity (4D GR) is approximately recovered at a scale below \( r_c \). However, the simple linear perturbation does not work in this regime. Only recently has the mechanism to recover 4D GR been clarified under the restriction to spherically symmetric configurations, and the leading correction to 4D GR has been derived. Here, we develop an alternative formulation which can handle more general perturbations. We also generalize the model by adding a bulk cosmological constant and the brane tension.

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I. INTRODUCTION

A braneworld model, whose gravity behaves as four dimensional at a short distance scale but shows a higher dimensional nature at larger distances, was proposed by Dvali, Gabadadze, and Porrati (DGP) [1]. In this model, the brane, on which the fields of the standard model are confined, has an induced Einstein-Hilbert term [2,3]. This model has various cosmologically interesting features [4–11]. Particularly in the model with a five-dimensional bulk, an interesting cosmological solution was found, in which the accelerated expansion of the Universe at a late epoch is realized without introducing the cosmological constant [7]. Based on this model, a novel mechanism that dilutes the cosmological constant was also proposed [11].

Although we mentioned above that gravity in this model at short distances is expected to behave as four dimensional, it is not so transparent if the model actually mimics four-dimensional general relativity (4D GR). The linear analysis of this model shows that the tensor structure of the induced metric perturbations takes a five-dimensional form even at short distance [1]. The situation is analogous to the case of models with massive gravitons. In this case the deviation from 4D GR does not vanish even in the massless limit, which is known as the van Dam–Veltman–Zakharov discontinuity [12–14]. In this context, the possibility that the 4D GR is recovered by nonlinear effect was suggested in Ref. [15]. There have been many discussions about this issue [16]. In particular, we have a clear statement that the discontinuity disappears when we introduce the cosmological constant [17,18]. Although the analysis with a cosmological constant is quite suggestive, the discontinuity is absent only when the limit is taken, keeping the length scale determined by the cosmological constant much smaller than the Compton wavelength of the massive graviton. In this paper, we develop an alternative formalism which can handle general perturbations in a weak gravity regime.

II. SETUP

The model that we consider is defined by the five-dimensional action

\[
S = \frac{M_5^2}{4r_c} \int d^5x \sqrt{-g} \left( R^{(5)} + \frac{12}{\ell^2} \right) + \int d^4x \sqrt{-g^{(4)}} \left( \frac{M_5^2}{2} R^{(4)} - \frac{3M_5^2}{r_c\ell} + L_{\text{matter}} \right),
\]

where \( M_5 \) is the five-dimensional graviton mass, \( \ell \) is the extra dimension scale, and \( L_{\text{matter}} \) is the matter content on the brane.
where $M_4$, $r_0$, and $\ell$ are constants. $R^{(5)}$ and $R^{(4)}$ are, respectively, the curvature scalars corresponding to the five-dimensional metric $g_{\mu\nu}$ and the four-dimensional one $g^{(4)}_{\mu\nu}$ induced on the brane. Here, we added both the bulk cosmological constant and the brane tension terms to the original DGP model. They are tuned to admit the Minkowski brane as a vacuum solution. The model is reduced to the original one by setting $\xi \to \infty$. The unperturbed background geometry is given by five-dimensional anti-de Sitter space-time,

$$
d s^2 = g^{(0)}_{ab} dx^a dx^b = dy^2 + \gamma_{\mu\nu}(y) dx^\mu dx^\nu
$$

with a brane located at $y=0$, where a $Z_2$ symmetry is imposed. Here $\eta_{\mu\nu}$ is a four-dimensional Minkowski metric.

### III. SEMINONLINEAR PERTURBATIONS

We follow the method of Ref. [29] introduced for the purpose of analyzing weak gravity in the Randall-Sundrum model [30]. We prepare two coordinate systems. In the coordinates $\{x^a\}$, the gauge is chosen so that the metric perturbations $h_{ab}$ can be easily computed in the five-dimensional bulk. That is, we use the Randall-Sundrum gauge,

$$
h_{5a} = 0, \quad h_{\mu} = 0, \quad h_{\mu\nu} = 0.
$$

In this paper the fifth direction is the direction of extra dimension. The Greek and Latin indices represent four- and five-dimensional coordinates, respectively. The other coordinate system $\{\bar{x}^a\}$ satisfies the Gaussian normal conditions

$$
\bar{h}_{5a} = 0,
$$

and also keeps the location of the brane unperturbed at $\bar{y} = 0$. Under the coordinate transformation $x^a = \bar{x}^a - \xi^a(\bar{x})$, the metric perturbation transforms as

$$
\begin{align*}
\bar{h}_{ab} &= h_{ab}[\bar{x} - \xi(\bar{x})] + \left(-g^{(0)}_{ab} \xi^2 + \frac{1}{2} g^{(0)}_{ab55} (\xi^2)^2 - \cdots \right) \\
&- \{\xi_a[g^{(0)}_{cb}(\bar{x} - \xi) + h_{cb}(\bar{x} - \xi)] + (a \leftrightarrow b) \} \\
&+ [g^{(0)}_{cd}(\bar{x} - \xi) + h_{cd}(\bar{x} - \xi)] \xi_{a} \xi_{b}.
\end{align*}
$$

The argument of the variables is supposed to be $\bar{x}$ unless otherwise is specified, and "_a_" denotes a differentiation with respect to $\bar{x}^a$.

The conditions that the $\{00\}$ component and $\{0\mu\}$ components are zero in both coordinates provide equations for the gauge parameters, which are solved up to second order as

$$
\begin{align*}
\xi^5 &= \xi^{(2)} + \xi^5, \\
\xi^\mu &= \xi^{(2)} (\eta^{\mu\nu} - \gamma^{\mu\nu}) \xi^5 + \xi^\mu + \xi^{(2)}.
\end{align*}
$$

where $\xi^5(\bar{x}^\nu)$ and $\xi^{(2)}(\bar{x}^\nu)$ are the values of the gauge parameters evaluated on the brane, and

$$
\begin{align*}
\xi^{(2)} &= \int_0^y d\bar{y} \gamma^{\mu\nu} \xi^{(2)}(\bar{x}^\nu) \xi^5 + \xi^\nu, \\
\xi^{(2)} &= \int_0^y d\bar{y} \gamma^{\mu\nu} \left[ \xi^{(2)}(\bar{x}^\nu + 2 \xi^\nu) \xi^5 - \xi^{(2)}(\bar{x}^\nu + \xi^\nu) \xi^5 - \xi^{(2)}(\bar{x}^\nu) \xi^5 \right].
\end{align*}
$$

We assume the following order counting:

$$
\begin{align*}
\xi^{(2)} &\ll \xi^5, \\
\xi^5 &\ll \frac{\xi^5}{\ell}, \\
\xi^{(2)} &\ll \frac{\xi^{(2)}}{\ell},
\end{align*}
$$

and keep the terms up to $O(\ell^2)$. Here $\ell^2$ is the order of the Newton potential $\Phi = -\frac{1}{2} h_{00}$. Later we will verify the consistency of this order counting. Then, the transformation for $\{\mu\nu\}$ components reduces to

$$
\begin{align*}
\bar{h}_{\mu\nu}(\bar{x}) &= h_{\mu\nu}(x - \xi(\bar{x})) + \delta h_{\mu\nu}, \\
\delta h_{\mu\nu} &= 2 \gamma^{\mu\nu} \xi^5 - \xi_{\mu\nu} - \xi_{\mu\nu} + \xi^5_{\mu\nu}.
\end{align*}
$$

Hereafter, the Greek indices are lowered or raised by the metric $\gamma_{\mu\nu}$.

The brane location is given by $\bar{y} = 0$. Hence in the $\{x^a\}$ coordinates the brane is bent. For simplicity, we impose the harmonic gauge condition, $\bar{h}^{\nu}_{\mu\nu} = \frac{1}{2} \bar{h}_{\mu\nu}$ for the induced metric on the brane. To second order in $\xi^5$, this condition gives

$$
\xi^5 = \delta^{-1} \left[ -2 \xi^5_{\mu\nu} + (\Box \xi^5) \xi^5_{\mu\nu} - \frac{4}{\ell^2} \xi^5_{\mu\nu} \right].
$$

From this relation, we find that the assumption $\xi^5_{\mu\nu} = O(\ell^5)$ is consistent if the assumed order of $\xi^5$ is correct. Substituting Eq. (3.9), the gauge transformation $\delta h_{\mu\nu}$ evaluated on the brane becomes

$$
\delta h_{\mu\nu}[y=0] = \frac{\xi^5}{\ell} \xi^5_{\mu\nu} + \delta^{-1} \left[ \frac{4}{\ell^2} \xi^5_{\mu\nu} + 2 \gamma^{\mu\nu} \xi^5_{\mu\nu} \xi^5_{\mu\nu} - 2 (\Box \xi^5) \xi^5_{\mu\nu} \right].
$$

Then, the trace of the induced metric is also evaluated as

$$
\bar{h} = \frac{12}{\ell} \xi^5 + 2 \delta^{-1} \left[ \gamma^{\mu\nu} \gamma^{\mu\nu} \xi^5_{\mu\nu} \xi^5_{\mu\nu} - (\Box \xi^5)^2 \right].
$$

Next we consider the junction condition. After a straightforward calculation, we can show
(\partial_x^2 + 2 \ell^{-1})\bar{h}_{\mu\nu} = (\partial_x^2 + 2 \ell^{-1})h_{\mu\nu}(\bar{x} - \xi) + 2 \xi^5
+ \frac{2}{\ell} \Xi_{\mu\nu} \quad (\text{at } \bar{y} = 0),
(3.12)

with

\Xi_{\mu\nu} := \gamma_{\mu\nu} \gamma^{\rho\sigma} \xi_{\nu,\sigma} + \xi_{\mu,\nu}.
(3.13)

Using this relation, the junction condition becomes

\begin{equation}
T_{\mu\nu} - \frac{1}{2} \gamma_{\mu\nu} T + \frac{M_2^2}{2} \Box \bar{h}_{\mu\nu} = - \frac{M_2^2}{2} \left[(\partial_x^2 + 2 \ell^{-1})h_{\mu\nu}(\bar{x} - \xi)\big|_{\bar{y}=0} + 2 \xi^5\right] \bar{h}_{\mu\nu}
+ \frac{2 \Xi_{\mu\nu}}{\ell} + \gamma_{\mu\nu} \left[\Box \xi^5 + \frac{\Xi}{\ell}\right].
(3.14)
\end{equation}

The equation that determines the brane bending is obtained from the trace of the above equation as

\begin{equation}
\Box^{-1} \frac{1}{M_4^2} T = \bar{h} + \frac{3}{r_c} \left( \xi^5 + \Box^{-1} \frac{\Xi}{\ell} \right)
= \frac{3}{r_c} \xi^5 + \Box^{-1} \left( \gamma_{\mu\nu} \gamma^{\rho\sigma} \xi^5 \xi_{\mu,\rho} \xi_{\nu,\sigma} - (\Box \xi^5)^2 \right)
+ \frac{3}{r_c} \frac{\Xi}{\ell} \Box^{-1} \Xi,
(3.15)
\end{equation}

where we have introduced \( r_c^p := [(1/r_c) + (2/\ell)]^{-1} \). We can neglect the last term on the right hand side of Eq. (3.15) because it is always higher order compared with the first term. The left hand side is something like the Newton potential \( \Phi \), hence we assume it to be \( O(\ell^2) \). Outside the matter distribution with the total mass \( m \), the left hand side can be expressed as \( \sim r_c/\ell \), where \( r_c := m 4\pi M_4^2 \). At large scale, \( r \cong (r_c r_c)^{1/3} \), the first term on the right hand side dominates, while at small scale, \( r \cong (r_c r_c)^{1/3} \), the second term dominates. Therefore we have

\begin{equation}
O(\ell^2) = \frac{3}{r_c} \xi^5 + \Box^{-1} \frac{\Xi}{\ell} \Xi.
(3.16)
\end{equation}

Thus we find that our assumption as to the order counting for \( \xi^5 \) is justified.

IV. MECHANISM FOR RECOVERING FOUR-DIMENSIONAL GENERAL RELATIVITY

The remaining task is to evaluate \( (\partial_x^2 + 2 \ell^{-1})h_{\mu\nu}(\bar{x} - \xi)\big|_{\bar{y}=0} \). Here we need to solve the bulk field equations. Different from the R-S case, we solve the bulk equations with the Dirichlet boundary condition (3.7). Here, we note that the location of the brane is not a straight sheet in the coordinates in the R-S gauge \( \{x^a\} \).

We can give the general solution for the bulk field equations as a superposition of homogeneous mode solutions with a purely outgoing boundary condition:

\begin{equation}
h_{\mu\nu}(x) = \int \mathcal{H}_{\mu\nu}(p)e^{ip_{\mu}x^\mu} K_2(p\ell e^\gamma/\ell) d^4 p
= \int \mathcal{H}_{\mu\nu}(p)e^{ip_{\mu}(\bar{\xi} - \xi)} K_2(p\ell e^\gamma - e^\gamma/\ell) d^4 p.
(4.1)
\end{equation}

where \( K_2(p\ell) \) is the modified Bessel function and \( \mathcal{H}_{\mu\nu} \) is the expansion coefficient. The coefficient \( \mathcal{H}_{\mu\nu} \) is to be determined so as to satisfy the Dirichlet boundary condition

\begin{equation}
h_{\mu\nu}|_{\bar{y}=0}(\bar{x}) = \int \mathcal{H}_{\mu\nu}(p)e^{ip_{\mu}(\bar{\xi} - \xi)} K_2(p\ell e^{-\gamma/\ell}) d^4 p.
(4.2)
\end{equation}

If we are allowed to approximate the above expression by setting \( \xi^a = 0 \), we have \( \bar{h}_{\mu\nu}(p) := (2\pi)^{-4} \int d^4 \bar{x} e^{-ip_{\mu}x^\mu} \times h_{\mu\nu}(\bar{x}) = \mathcal{H}_{\mu\nu}(p) \), and therefore we have

\begin{equation}
(2\pi)^{-4} \int d^4 \bar{x} e^{-ip_{\mu}x^\mu}(\partial_x^2 + 2 \ell^{-1})h_{\mu\nu}|_{\bar{y}=0}(\bar{x})
= - \frac{pK_1(p\ell)}{K_2(p\ell)} \bar{h}_{\mu\nu}(p)
- \frac{pK_1(p\ell)}{K_2(p\ell)}(\bar{h}_{\mu\nu} - \delta \bar{h}_{\mu\nu}).
(4.3)
\end{equation}

We think that the errors caused by this naive approximation are not large, although any rigorous proof is not ready yet. If the leading errors are simply proportional to \( \bar{h}_{\mu\nu} \), we can neglect them since they are of higher order in \( \epsilon \). Such a naive expansion with respect to \( \xi \) will be justified for small \( p \). But for large \( p \), we will not be allowed to expand the combination \( p \xi \) in the exponent. However, as we will see below, even the leading correction to the gravitational potential coming from the contribution of this part is suppressed to be irrerelevantly small at small scale \( r \ll r_c \). Hence, the errors due to this naive approximation can be crucial only if this approximation significantly underestimate the magnitude of \( (\partial_x^2 + 2 \ell^{-1})h_{\mu\nu}|_{\bar{y}=0} \), which is quite unlikely.

Using Eq. (4.3), the junction condition (3.14) is written down explicitly as

\begin{equation}
\bar{h}_{\mu\nu} - \frac{2}{M_4^2} \left( \bar{T}_{\mu\nu} - \frac{1}{2} \gamma_{\mu\nu} \bar{T} \right)
= \sum \left[ -\frac{\Sigma_{\mu\nu}}{r_c} + \frac{2pK_1(p\ell)}{r_c K_2(p\ell)} \gamma_{\mu\nu} \xi^5 + \frac{2}{r_c} p \mu D^5 \xi^5
+ \frac{pK_1(p\ell)}{r_c K_2(p\ell)} \mathcal{D} \bar{h}_{\mu\nu} \right]
\left( \frac{\Xi_{\mu\nu}}{2} + \frac{1}{2} \gamma_{\mu\nu} \Xi \right)
(4.4)
\end{equation}

(at \( \bar{y} = 0 \),

where
deviation from 4D GR at small scale is less strongly sup-
critically we take the limit
and

Here the quantity with \( \tilde{ } \) represents the Fourier coefficient of the corresponding variable as before. We show that in the square brackets on the right hand side of Eq. (4.4) the first term gives the dominant contribution. We can drop the last two terms simply because they are always higher order in \( \epsilon \) compared with the first term. The second term is irrelevant since it can be eliminated by a four-dimensional gauge transformation. As a result, the equation that determines the metric induced on the brane \( \tilde{h}_{\mu \nu} \) is reduced to the one for the linear theory. The only difference is in the equation that determines the brane bending [Eq. (3.15)].

The order of magnitude of the first term on the right hand side of Eq. (3.15) is estimated as \( \xi^3 |r_c| \Phi^{-1/2} r/r_c \) at small scale, and hence it is suppressed by the factor \( \Phi^{-1/2} r/r_c \) compared with the Newton potential \( \Phi \). The leading behavior of the induced metric is therefore determined by setting the left hand side of Eq. (4.4) to zero. Thus we conclude that 4D GR is recovered by taking into account the nonlinear brane bending for weak gravity at small scale \( r \leq (r_c^2 r_g)^{1/3} \). If we take the limit \( r \rightarrow \infty \), all length scales come into this regime. Hence, we have confirmed the absence of van Dam–Veltman–Zakharov discontinuity. As first pointed out in Ref. [21], however, because of the factor \( \Phi^{-1/2} \) the leading order deviation from 4D GR at small scale is less strongly suppressed than the naively expected suppression \( r/r_c \).

At large scale, this term becomes more and more important. For \( r \geq (r_c^2 r_g)^{1/3} \), we have

\[
\tilde{\xi}^5_{\infty} = - \frac{r_c^2 \tilde{T}}{3 M_4^2 p^2}. \tag{4.7}
\]

This is nothing but the result for the linearized case. In Sec. V, we discuss the regime where the linear theory is valid. After that, in the succeeding section, we discuss the leading order correction to the 4D GR at short distance scale assuming static and spherically symmetric configurations.

V. LINEAR REGIME

In this section we consider perturbations at large scale \( r \gg (r_c^2 r_g)^{1/3} \), where the linear theory is valid. Substituting Eq. (4.7) into Eq. (4.4), we obtain

\[
\hat{h}_{\mu \nu} = \frac{2}{M_4^2} \left( \tilde{T}_{\mu \nu} - \frac{1}{2} \alpha \gamma_{\mu \nu} \tilde{T} \right), \tag{5.1}
\]

with

\[
D = \frac{1}{p^2 + \frac{p K_1(p \ell)}{r_c K_2(p \ell)}} \tag{4.5}
\]

and

\[
\delta h^{(2)}_{\mu \nu} = \xi^{-1} \left[ 2 \gamma^\mu \xi^5 \gamma^\nu,_{\rho \sigma} - 2 \left( \square \xi^5 \right) \delta^\mu_{\rho} \delta^\nu_{\sigma} \right]. \tag{4.6}
\]

Here the quantity with \( \hat{ } \) represents the Fourier coefficient of the corresponding variable as before. We show that in the square brackets on the right hand side of Eq. (4.4) the first term gives the dominant contribution. We can drop the last two terms simply because they are always higher order in \( \epsilon \) compared with the first term. The second term is irrelevant since it can be eliminated by a four-dimensional gauge transformation. As a result, the equation that determines the metric induced on the brane \( \tilde{h}_{\mu \nu} \) is reduced to the one for the linear theory. The only difference is in the equation that determines the brane bending [Eq. (3.15)].

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First we look at the behavior of the propagator \( D \), which was already discussed in Ref. [28]. When we consider the length scale much smaller than \( \ell \), we have

\[
D \rightarrow \frac{1}{p^2 + \frac{r_c^{-1} p}{r_c}} \quad (p \ell \gg 1). \tag{5.3}
\]

At a length scale smaller than \( r_c (p \ell \gg 1) \), the propagator \( D \) behaves as that for a five-dimensional field. When the length scale is much larger than \( \ell \), \( K_1(p \ell)/K_2(p \ell) \) goes to \( p \ell/2 \). Thus we have

\[
D \rightarrow \frac{2 r_c \ell}{(\ell + 2 r_c) p^2} \quad (p \ell \ll 1). \tag{5.4}
\]

Hence again the propagator \( D \) behaves as a four-dimensional field, but Newton’s constant is not given by \( 2 M_4^2 \) but by \( 2 M_4^{-2}/(1 + \ell/2 r_c) \).

Next we turn to the tensor structure specified by \( \alpha \). For a four-dimensional massless graviton we have \( \alpha = 1 \), while \( \alpha = \frac{1}{2} \) for the case of a massive graviton. For \( p \ell \gg 1 \), we have

\[
\alpha \rightarrow \frac{1 + \ell}{3 r_c}. \tag{5.5}
\]

We have \( \alpha \rightarrow 1 \) for \( \ell \gg r_c \), while \( \alpha \rightarrow 2/3 \) for \( \ell \ll r_c \). On the other hand, for \( p \ell \ll 1 \), we have \( \alpha \rightarrow 1 \) irrespective of the ratio between \( \ell \) and \( r_c \).

The results are summarized in Fig. 1. When \( r \gg \ell \), the 4D GR is realized by the Randall-Sundrum mechanism. The effective Planck mass differs from \( M_4 \) in this case. On the other hand, when \( r \ll r_c \), the gravity becomes four dimensional again, but the tensor structure differs from 4D GR.
VI. STATIC SPHERICAL SYMMETRIC NONRELATIVISTIC STAR

We consider a static spherical symmetric nonrelativistic star. Assuming that the energy momentum tensor is dominated by the \( \{00\} \) component \( T_{00}=\rho \), we neglect the effect of pressure when we solve the metric perturbation. First we solve the nonlinear equation for \( \xi^5 \) [Eq. (3.15)]. Under the present assumptions, Eq. (3.15) is simplified as

\[
-\frac{1}{M_4^2} \rho = \frac{1}{r^2} \partial_r \left( \frac{3r^2}{r_c^3} \xi^5 - \frac{2}{M_4^2} \frac{d}{dr} \left( \frac{\xi^5 r^2}{r_c^3} \right)^2 \right). \tag{6.1}
\]

This equation can be immediately integrated once, and we obtain

\[
2 \left( \xi^5_r \right)^2 - \frac{3r}{r_c^3} \xi^5_r - \frac{r_g(r)}{r} = 0, \tag{6.2}
\]

where

\[
r_g(r) = \frac{1}{M_4^2} \int_0^r drr^2 \rho. \tag{6.3}
\]

Outside the star, we have \( r_g(r)=\rho_g=ml/4\pi M_4^2 \). Solving the above equation with respect to \( \xi^5_r \), we have

\[
\xi^5 = \int dr \left[ \frac{3r}{4r_c^3} - \frac{1}{4} \sqrt{\left( \frac{3r}{r_c^3} \right)^2 - \frac{8r_g(r)}{r}} \right]. \tag{6.4}
\]

Here we have chosen the signature in front of the square root imposing the condition that \( \xi^5_r \) does not become large at \( r \to \infty \). The other branch with the “+” sign is outside the scope of the present formalism since we have assumed a Minkowski brane background. At small scale, this expression reduces to \( \xi^5 = -\int dr \sqrt{r_g(r)/2r} \). Outside the matter distribution, we simply have \( \xi^5 = -\sqrt{2r_g} \). Hence, the correction to the Newton potential is given by

\[
\delta \Phi \approx \sqrt{\frac{1}{2r}} \left( \frac{r}{r_c} \right)^2. \tag{6.5}
\]

which recovers the result obtained in Ref. [21].

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