Cosmic Microwave Radiation Anisotropies in Brane Worlds

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We propose a new formulation to calculate the Cosmic Microwave Background (CMB) spectrum in the Randall Sundrum two-branes model based on recent progresses in solving the bulk geometry using a low energy approximation. The evolution of the anisotropic stress imprinted on the brane by the 5D Weyl tensor is calculated. An impact of the dark radiation perturbation on CMB spectrum is investigated in a simple model assuming an initially scale-invariant adiabatic perturbations. The dark radiation perturbation induces isocurvature perturbations, but the resultant spectrum can be quite different from the prediction of simple mixtures of adiabatic and isocurvature perturbations due to Weyl anisotropic stress.

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1. Introduction

By suggestions from string theory/M-theory, much attention is paid to brane world ideas where we are living on a 3-brane in higher-dimensional spacetime. In these brane world models, only gravity propagates in a higher-dimensional "bulk" spacetime while standard model fields are confined to the brane. The simplest realization of this idea was given by Randall and Sundrum \[1\]. The 5D bulk has a negative cosmological constant \(\Lambda_5 = -6/l^2\) where \(l\) is the curvature radius in the bulk and the brane has a tension \(\lambda = 6/k_{\text{5D}}^2\) where \(k_{\text{5D}}^2 = 8\pi G_5\) and \(G_5\) is the 5D Planck constant.

In this model, predictions of 4D general relativity are modified due to the influence of the gravitational fields in the bulk. These modifications need to be consistent with cosmological observations. Among them, the observations on Cosmic Microwave Background (CMB) anisotropies are now dramatically improving. The WMAP experiment has already provided precise measurements of CMB anisotropies \[2\]. It is then necessary to calculate the CMB spectrum in brane worlds in order to check the consistency of the model.

These works were initiated soon after the paper of Randall Sundrum \[3\]. Although a large number of papers have tried to give a prediction of CMB anisotropy, there is still no quantitative estimation. This is due to the difficulties in solving full 5D perturbations \[4\]. In view of observational improvements, it is eagerly desired to develop a formulation which enables us to provide quantitative predictions. In this Letter, we propose such a formulation based on recent progresses in solving bulk geometry using a low energy approximation \[5\] \[6\] \[7\] \[8\].

We consider a simple Randall-Sundrum type two-branes model where a low energy approximation is applicable. The stabilization mechanism is not introduced and the physical brane is assumed to be the positive tension brane. This simple setup enables us to understand clearly how the bulk gravitational fields affect the evolution of perturbations on the brane.

2. View from the brane

The effective 4D Einstein equation on the brane is given by \[9\]:

\[
G_{\mu\nu} = \kappa_4^2 T_{\mu\nu} + \kappa_4^2 \Pi_{\mu\nu} - \epsilon_{\mu\nu},
\]

where \(\kappa_4^2 = 8\pi G_4 = \kappa_5^2/l^2\) and \(\Pi_{\mu\nu}\) is the quadratic function of the energy momentum tensor, which can be neglected at low energies \(T_{\mu\nu}/\lambda \ll 1\). The energy-momentum tensor satisfies the conservation’s law \(\Pi_{\mu\nu} = 0\) on the brane \(\text{Eq.(2)}\). One needs 5D equations for \(E_{\mu\nu}\) in the bulk and this is the source of great complexity of the problem. To observe this fact, it is convenient to parameterize \(\epsilon_{\mu\nu}\) as an effective energy-momentum tensor:

\[
\epsilon_{\mu\nu} = \kappa_4^2 \begin{pmatrix} -\left(\rho_{\epsilon} + \delta \rho_{\epsilon} Y\right) & aV_{\epsilon} Y_i \\ -a^{-1}V_{\epsilon} Y_j & \left(P_{\epsilon} + \delta P_{\epsilon}\right)\delta_{ij} + \delta\pi_{\epsilon} Y_{ij} \end{pmatrix},
\]

where \(Y(k, x) \propto e^{ikx}\) is the normalized scalar harmonics and the vector \(Y_i\) and traceless tensor \(Y_{ij}\) are constructed from \(Y\) as \(Y_i = -k^{-1} Y_{,i}\), \(Y_{ij} = k^{-2} Y_{,ij} + \delta_{ij} Y/3\).

In the background universe, \(\text{Eq.(2)}\) gives \(\rho_{\epsilon} + 4H\rho_{\epsilon} = 0\) and \(P_{\epsilon} = \rho_{\epsilon}/3\), where \(H = \dot{a}/a\) and \(a\) is the scale factor. The solution for Weyl energy density is then given by \(\rho_{\epsilon} = C a^{-4}\) where \(C\) is the integration constant. Thus the contribution from \(\epsilon_{\mu\nu}\) appears as dark radiation. It has been shown that \(C\) is related to the mass of the AdS-Schwarzschild BH in the bulk. In the following, we will assume \(\rho_{\epsilon} = 0\) in the background spacetime.
Let us consider the perturbations around the above background spacetime. The linear scalar metric is taken as

\[ ds^2 = -(1 + 2\Psi(t)Y)dt^2 + a(t)^2(1 + 2\Phi(t)Y)\delta_{ij}dx^idx^j. \]

Eq. (2) give

\[
\dot{\rho}_\xi + 4H\rho_\xi = -a^{-1}kV_\xi, \quad \dot{P}_\xi = \frac{1}{3}\delta\rho_\xi,
\]

\[
\dot{\xi}_\xi + 4HV_\xi = \frac{ka^{-1}}{3}(\delta\rho_\xi - 2\delta\xi_\xi), \quad (3)
\]

for \(\rho/\lambda \ll 1\). At large scales \(ka^{-1}/H \to 0\), we can neglect the terms proportional to \(k\). The solution for \(\delta\rho_\xi\) is given by \(\delta\rho_\xi = \delta Ca^{-4}\) where \(\delta C\) is the integration constant. The amplitude \(\delta C\) is related to the perturbatively small AdS-Schwarzschild mass in the bulk. A problem is that it is impossible to determine anisotropic stress \(\delta\xi_\xi\) because it is dropped from Eq. (3) for \(ka^{-1}/H \to 0\). In [11], it was clearly shown that this uncertainty of \(\delta\xi_\xi\) prevents us from predicting CMB anisotropy. The temperature anisotropy caused by (direct) Sachs-Wolfe effect is given by the formula [11]

\[
\frac{\Delta T}{T} = \Theta_0 + \Psi = \zeta + \Phi - \Phi, \quad (4)
\]

where \(\Theta_0\) is the temperature anisotropy of radiation and we used Eq. (1) in the second equality. The evolution of \(\zeta = \Phi + \delta\rho/3(\rho + P)\) is determined only by the conservation of the matter energy-momentum tensor. We will assume the matter perturbation is adiabatic \(\delta P = c_s^2\delta\rho\) where \(\delta\rho\) is the density perturbation, \(\delta P\) is the pressure perturbation and \(c_s\) is the sound velocity. The continuity equation for matter perturbation implies \(\zeta = \zeta_s = \text{const}\). In addition to \(\zeta\), the solutions for metric perturbations \(\Phi\) and \(\Psi\) are needed. From the effective 4D equation [11], we can get the equations to determine \(\Phi\) and \(\Psi\) as

\[
\zeta_{tot} = \Phi - \frac{H}{\dot{H}}\left(\frac{\dot{\Phi}}{H} - \dot{\Psi}\right), \quad \Phi + \Psi = -\kappa_2^2k^{-2}a^2\delta\xi_\xi, \quad (5)
\]

where \(\zeta_{tot}\) is the curvature perturbation on hypersurface of uniform total energy density;

\[
\zeta_{tot} = \zeta + \frac{\delta\rho_\xi}{3(\rho + P)} = \zeta_s + \frac{\delta Ca^{-4}}{3(\rho + P)}. \quad (6)
\]

Note that the solution for \(\zeta_{tot}\) was obtained only by the law of conservation. However, the latter equation in Eq. (3) contains undetermined \(\delta\xi_\xi\). Indeed, Eq. (3) is written at the decaying time as

\[
\frac{\Delta T}{T} = -\frac{1}{3}\zeta_s - \frac{2}{3}\left(\frac{\rho_\xi}{\rho}\right)\delta C_s
\]

\[
- \kappa_2^2k^{-2}a^2\delta\xi_\xi + 2\kappa_2^2k^{-2}a^{-5/2}\int a^{7/2}\delta\xi_\xi da_s, \quad (7)
\]

where \(\delta C_s = \delta\rho_\xi/\rho_s\) and \(\rho_s\) is the radiation energy density. The second term comes from the entropy perturbation contribution due to Weyl energy density. Unless the behavior of \(\delta\xi_\xi\) is known, we cannot say anything about the effect of this dark radiation. It is possible to assume some ansatz for \(\delta\xi_\xi\), say \(\delta\xi_\xi \propto \delta\rho_\xi\) [11]. However, because \(\delta\xi_\xi\) is induced by 5D gravitational perturbations, it is essential to solve the evolution equations for \(\xi_{\mu\nu}\) in the bulk.

3. Solution for \(\xi_{\mu\nu}\) in the bulk

The evolution equations for \(\xi_{\mu\nu}\) obtained from 5D Bianchi identity are consisted of discouragingly complicated partial differential equations. Fortunately, the curvature radius on the brane defined by \(L^2 \sim 1/\nu_0^2 \sim 1/\kappa_0^2T_0^\mu\) at the decaying time is significantly longer than the curvature radius in the bulk \(L\) because the experiments on Newton’s force law already impose the constraint on \(l < 1\) mm. Thus the system has a natural small parameter \(\epsilon = (l/L)^2 \sim T_0^\mu/\lambda\) and it is possible to solve the equations by the perturbation in terms of \(\epsilon\). If we consider a second brane with negative tension \(\lambda_c = -6/\kappa_0^2l\) at the physical distance \(d_0(x)/l\) from our brane, \(\xi_{\mu\nu}\) can be determined as [7, 8]

\[
\xi_{\mu\nu}^\mu = \frac{2}{\epsilon^{2d_0} - 1} \left[ -\frac{\kappa_0^2}{2}(T_0^\mu + e^{-2d_0T_0^\mu}) - \nabla^\mu\nabla_\nu d_0 + \delta_{\nu}^\mu\nabla^2 d_0 - \left(\nabla^\mu d_0\nabla_\nu d_0 + \frac{1}{2}\delta_{\nu}^\mu(\nabla d_0)^2\right) \right], \quad (8)
\]

where \(T_0^\mu\) is the energy-momentum tensor on the second brane. Because the bulk spacetime shrinks exponentially due to the negative cosmological constant in the bulk, the curvature radius on the second brane \(L_c\) is shorter for larger \(d_0\) as \(L_c = Le^{-d_0}\). Thus in order to ensure that the low energy approximation can be applied on the second brane \((1/L_c)^2 < 1\), the radion \(d_0\) should satisfy \(e^{-2d_0} > (1/L)^2 \ll 1\).

The evolution of \(d_0\) can be calculated from the traceless condition \(\xi_{\mu\nu}^\mu = 0\) in Eq. (2):

\[
\nabla^2 d_0 - (\nabla d_0)^2 = \frac{\kappa_0^2}{6}(T + e^{-2d_0T_0}). \quad (9)
\]

From Eqs. (8) and (9) the behavior of \(\xi_{\mu\nu}\) is completely determined. Substituting the solution for \(\xi_{\mu\nu}\) into Eq. (1), the effective theory on the brane becomes quasi-scalar-tensor theory [12].

4. View from the bulk

The solution for \(\xi_{\mu\nu}\) (Eq. (8)) should be consistent with Eq. (2). In the background spacetime, \(\rho_\xi\) is written in terms of \(d_0\) and energy densities on both branes;

\[
\kappa_2^2\rho_\xi = \frac{1}{\epsilon^{2d_0} - 1}\left(6Hd_0 - 3d_0^2 - \kappa_2^2(\rho + e^{-2d_0}\rho_s)\right).
\]

The evolution equation for \(d_0\) can be integrated once with the help of the energy-momentum conservation on each
brane. We get $\rho_c = Ca^{-4}$ where $C$ is the integration constant. In this case, the integration constant is interpreted as the initial condition for the radion $d_0(t)$. Now we turn to the perturbations. We denote the perturbation of the radion as $d_0(t,x) = d_0(t) + N(t)Y(k,x)$. $\delta \rho_c$ can be evaluated in the same way as the background spacetime. We get $\delta \rho_c = \delta Ca^{-4}$ at large scales where $\delta C$ is the integration constant associated with $N$. An advantage of our approach is that an equation for Weyl anisotropic stress $\delta \pi_\xi$ can be derived as

$$\kappa_4^2 a^{-2} \delta \pi_\xi = \frac{2}{e^{2d_0} - 1} N.$$  \hfill (10)

Hence the behavior of $\delta \pi_\xi$ is determined by the radion perturbation $N$. Because the radion perturbation $N$ is coupled to the metric perturbations $\Phi$ and $\Psi$, we should solve them at the same time. The behavior of the metric perturbations is determined by imposing the equation of state on the matter perturbations: $\delta P = c_s^2 \delta \rho$ and $\delta P_\xi = c_s^2 \delta \rho_c$. We then have three equations, $\delta P = c_s^2 \delta \rho$, $\delta P_\xi = c_s^2 \delta \rho_c$, and $\Phi + \Psi = -\kappa_4^2 a^{-2} \delta \pi_\xi$ for three unknown functions, $N, \Psi$ and $\Phi$. It is convenient to introduce a new set of variables $\varphi, \varphi_c$ and $E$ to solve these equations. We take

$$\Phi = -\varphi - k^{-2} a^2 H \dot{E} + \frac{1}{3} E, \quad N = \varphi_c - \varphi - k^{-2} a^2 d_0 \dot{E},$$

$$\Psi = -\varphi - k^{-2} a^2 (\dot{E} + 2H \dot{E}).$$ \hfill (11)

The equation $\delta P = c_s^2 \delta \rho$ gives the equation for $\varphi$;

$$\ddot{\varphi} + (2 + 3c_s^2) \dot{H} \dot{\varphi} - (3H^2 + 2 \dot{H} + 3c_s^2 H^2) \varphi - \frac{1}{2} \left( \frac{1}{3} - c_s^2 \right) \kappa_4^2 \delta \rho_c \varphi = -k^2 a^{-2} \varphi \left[ \left( c_s^2 \frac{2}{3} \right) - \frac{1}{3} \left( c_s^2 + \frac{1}{3} \right) E \right].$$ \hfill (12)

At large scales $ka^{-1}/H \rightarrow 0$, the equation for $\varphi$ is decoupled from $E$ and it can be integrated once to give the first order differential equation for $\varphi$;

$$- \left[ \dot{\varphi} - \frac{H^2}{\dot{H}} \left( \dot{\varphi} \frac{\dot{H}}{H} - \varphi \right) \right] = \varphi_c + \frac{\delta Ca^{-4}}{3(\rho + P)}.$$ \hfill (13)

where $\varphi_c$ is the integration constant. We should note that the left-hand side of Eq. (13) is nothing but the solution for $\zeta_{\text{tot}}$ (see Eqs. (9) and (11)). Then the behavior of $\varphi$ determines the evolution of the curvature perturbation which is independent of $d_0$. The evolution equation for $\varphi_c$ is obtained from $\delta P_\xi = c_s^2 \delta \rho_c$ and it is given by replacing $H, a, c_s^2, d/dt$, and $\delta \rho_c$ to $H_e = e^{d_0}(H - \dot{d_0}), a_e = ae^{-d_0}, c_{s_e}^2, e^{d_0}d/dt$ and $e^{d_0}\delta \rho_e$, respectively, in Eq. (12). The function $\varphi$ and $\varphi_c$ describe the displacement of our brane and the second brane respectively, and their relative difference causes the radion perturbation $N$. The equation $\Phi + \Psi = -\kappa_4^2 a^{-2} \delta \pi_\xi$ gives the evolution equation for $E$;

$$\ddot{E} + \left( 3H + \frac{2\dot{d_0}}{e^{2d_0} - 1} \right) \dot{E} - \frac{1}{3} k^2 a^{-2} E = \frac{2e^{2d_0}}{e^{2d_0} - 1} k^2 a^{-2}(\varphi - e^{-2d_0} \varphi_c).$$ \hfill (14)

The function $E$ is identified with the bulk anisotropic perturbation $E$. We have a closed set of equations for $\varphi, \varphi_c$ and $E$. Therefore, the solution for $\delta \pi_\xi$ can be obtained from Eqs. (10) and (11).

5. CMB anisotropy in a simple model

Let us consider the simplest case where $d_0 = d_* = \text{const.}$ and $d_*$ is sufficiently large $e^{-2d_*} \ll 1$. It can be realized by taking $\rho_c = -\rho e^{2d_*}$ and $w = w_c$. This choice is consistent with $\rho_e = 0$. First, let us consider large scale perturbations. Assuming the scale factor is given by $a \propto t^{2/3(1+w)}$ ($w =$ const.), the solution for $\varphi$ is obtained as

$$\varphi = -3(1 + w) \varphi_c - \frac{1}{9w - 1} \left( \frac{\rho_c}{\rho} \right) \delta C_\ast,$$

$$\varphi_c = -3(1 + w) \varphi_c - \frac{1}{9w - 1} \left( \frac{\rho_c}{\rho} \right) e^{2d_*} \delta C_\ast,$$

for $w \neq -1/3, 1/9$, where $\varphi_c$ is the curvature perturbation on hypersurface of uniform matter energy density on the second brane. Here we neglected the homogeneous solutions. Note that $w = 1/9$ is not singular and one can find a solution for $\varphi = -(1/2)\delta C_\ast(\rho/\rho_c)\ln a$. One point is that $\varphi_c$ that depends on $\delta C_\ast$ satisfies $\varphi_c = e^{2d_*}\varphi$, then $\delta C_\ast$ does not contribute to $E$. Now we can obtain the solution for metric perturbations. We find the parts of $\Phi$ and $\Psi$ corresponding to each term in the solution for $\varphi$:

$$\Psi = \Psi_\xi + \Psi_\xi, \quad \Phi = \Phi_\xi + \Phi_\xi, \quad \Psi_\xi = \Phi_\xi = -\varphi_c,$$

$$\varphi_c = -\frac{1}{9w - 1} \delta C_\ast \left( \frac{\rho_c}{\rho} \right).$$ \hfill (15)

Here, $\Phi_\xi$ and $\Psi_\xi$ are the same as conventional 4D solutions except for the additional terms which depend on the radion $d_*$ and "shadow matter" $\varphi_c$. These additional terms can be neglected for $e^{-2d_*} \ll 1$. Then we find the solution for Weyl anisotropic stress as

$$\kappa_4^2 a^{-2} \delta \pi_\xi = 2\varphi_c.$$ \hfill (16)

Large scale CMB anisotropy can be written as

$$\frac{\Delta T}{T} = -\frac{1}{5} \varphi_c.$$ \hfill (17)

Dark radiation does not affect $\Delta T/T$. This is quite a non-trivial result. As mentioned before, the Weyl energy density perturbations induce the entropy perturbations. The effect of Weyl anisotropic stress $\delta \pi_\xi$ exactly cancels this entropy perturbation in Eq. (17).
Let us investigate whether this cancellation holds for small scale perturbations. Under a tight coupling approximation of baryon-photon fluid, the evolution equation for radiation temperature anisotropy $\Theta_0$ becomes 

$$
\dot{\Theta}_0 + H\dot{\Theta}_0 + \frac{\dot{R}}{1+R} \dot{\Theta}_0 + k^2a^{-2}c_s^2\Theta_0 = F,
$$

$$
F = -\dot{\psi} - H\dot{\psi} - \frac{\dot{R}}{1+R} \dot{\psi} - \frac{k^2}{3}a^{-2}\Psi,
$$

(18)

where $R = 3\rho_b/4\rho_r$ and $\rho_b$ is the baryon energy density. The solutions for metric perturbations are given by the upper equation in Eq. (15) though $\varphi_\xi$ is the solutions for Eq. (12) with $E = 0$ and the equation for $\delta\varphi_E$ is obtained by Eqs. (13) and (14). The source term $F$ can be decomposed as $F = F_\xi + F_\xi$ according to the solutions for metric perturbations (Eq. (15)). Observed temperature anisotropy (Eq. 1) is given by $\Delta T/T = \Theta_0 + \Psi = \Theta_0 - \varphi_\xi + \Psi_\xi$. Thus if $\Theta_0 = \Theta_0 - \varphi_\xi$ behaves in the same way as conventional 4D theory, the temperature anisotropy also behaves in the same way. From Eq. (18), we get the equation for $\Theta_{0e}$:

$$
\ddot{\Theta}_{0e} + H\dot{\Theta}_{0e} + \frac{\dot{R}}{1+R} \dot{\Theta}_{0e} + k^2a^{-2}c_s^2\Theta_{0e} = F_\xi + \left(\frac{1}{3} - c_s^2\right)k^2a^{-2}\varphi_\xi.
$$

(19)

Hence at large scales $k \to 0$ or in the radiation dominated universe $c_s^2 = 1/3$, the temperature anisotropy is exactly the same as the conventional 4D theory. But as the Universe becomes matter dominated $c_s^2 \neq 1/3$ and perturbations enter the horizon, the behavior of temperature anisotropy is modified. Because the amplitude of $\varphi_\xi$ decreases with time under the horizon, the effect is largest for the first acoustic oscillation. Fig. 1 shows the resultant CMB anisotropy for given cosmological parameters. One can see that the dark radiation does not modify the CMB anisotropy at $l \sim 2$ and at large $l$ but it modifies a first acoustic oscillation. The resultant spectrum is quite different from the prediction of simple mixtures of adiabatic and isocurvature perturbations due to Weyl anisotropic stress $\delta\pi_\xi$.

6. Discussions

In this letter, we developed a formulation to calculate CMB anisotropy in two branes model at low energies with an appropriately small distance between two branes. In realistic models, we should introduce stabilization mechanisms of the radion. Our formulation can be extended to include stabilization mechanisms. We considered only the classical theory of the perturbations. Thus the initial spectrum remains to be determined, and it can be modified by the effects from the bulk in the brane world. These issues need further investigations.

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FIG. 1: CMB angular power spectrum for various $\delta C_\zeta$, $\zeta_*$ is appropriately normalized. We take $\Omega_0 = 0.26, \Omega_\Lambda = 0.70, \Omega_m = 0.04, h = 0.72$ and $n = 1$. The observational data was taken from WMAP [2] and the spectrum was calculated by modifying CMBEASY which is based on CMBFAST [14].