Gravitational Mass in Asymptotically de Sitter Space-Times with Compactified Dimensions

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We define gravitational mass in asymptotically de Sitter space-times with compactified dimension. It was shown that the mass can be negative for space-time with matter spreading beyond the cosmological horizon scale or large outward ‘momentum’ in four dimension. We give simple examples with negative energy in higher dimensions even if the matter is not beyond horizon or system does not have large ‘momentum’. They do not have the lower bound on the mass. We also give a positive energy argument in higher dimensions and realise that elementary fermion cannot exist in our examples.

I. INTRODUCTION

Superstring or M-Theory may offer the proper theory of gravity [1]. Such kind of theories are formulated in higher dimensions and it is believed that the extra dimension space will be compactified to be less than Planck length.

The stability of such space-times is important and has been discussed. It was shown that there is the instanton which may indicate the decay of Kaluza-Klein vacuum [2] [3] [4]. As Witten pointed out, the decay mode is excluded by the existence of (massless) elementary fermion related to supersymmetry [5].

On the other hand, the stability of the asymptotically anti-de Sitter(AdS) space-time with compactified dimension has been focused recently by Horowitz & Myers [6] because AdS/CFT correspondence [7] links the stability of super-Yang-Mills theory to that of AdS space-time. So they suggested a positive energy conjecture in locally asymptotically AdS space-times.

In the actual cosmological context, same argument of stability is important. In this paper, for simplicity, we consider D-dimensional Einstein gravity with a positive cosmological constant. A positive cosmological constant is essential for inflation universe [8] and acceleration of universe confirmed gradually by recent observations of supernovae [9]. In Nakao et al’s study, it was shown that the gravitational mass can be negative in four dimensional space-times when the matter distributes beyond the cosmological horizon scale [10] [11]. In other words, the mass is negative if the space-time has a large outward ‘momentum’. Evaluating the electric part of Weyl tensor, they also checked that the mass is related to gravitational tidal force. In this paper we find a initial data set whose mass can be negative and does not have the lower bound even if the space-time does not have ‘momentum’. Furthermore, we give an other dynamical solution which also does not have ‘momentum’.

In asymptotically flat space-times with compactified extra dimensions, there are examples of momentarily static initial slices such that the energy can be negative regardless of the size of the compactified dimensions [4]. A parallel

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argument helps us to discuss about the lower bound of the energy in asymptotically de Sitter space-times. Since a positive cosmological constant prompts the rapid expansion of the universe, we might be able to obtain an implication about the effect to the energy from the dynamics of the compactified dimensions.

The rest of the present paper is organised as follows. In the next section we give the definition of gravitational mass in asymptotically de Sitter space-time with extra dimensions and write down the expression in terms of canonical quantities on the hypersurfaces. Since a part of this is straightforward extension of work by Nakao et al [10] based on [12] [13], we give the brief discussion. We will point out that the expression is a special case of the gravitational Hamiltonian defined by Hawking & Horowitz [4]. In Sec. 3 we show two examples which have negative energy and no naked singularity. In Sec. 4, we consider the positive energy theorem to discuss the stability of de Sitter space-time. As a result we confirm that elementary fermion cannot exist in examples given by us. Finally we give summary in Sec. 4.

II. GRAVITATIONAL MASS IN ASYMPTOTICALLY DE SITTER SPACE-TIMES WITH HIGHER DIMENSIONS

We consider D-dimensional space-times which satisfy the Einstein equation with a positive cosmological constant $\Lambda$,

$$R_{IJ} - \frac{1}{2}g_{IJ}R + \Lambda g_{IJ} = 8\pi G_D T_{IJ},$$

(2.1)

where suffices $I, J$ runs over $0, 1, ..., D - 1$, $T_{IJ}$ is the energy-momentum tensor and $G_D$ is D-dimensional Newton’s constant. We decompose the space-time metric into the D-dimensional de Sitter metric $\bar{g}_{IJ}$ and the rest $h_{IJ}$:

$$g_{IJ} = \bar{g}_{IJ} + h_{IJ}.$$  

(2.2)

Hereafter the notation ‘over-bar’ indicates the quantities of the background de Sitter space-time. We remember that $h_{IJ}$ is not necessarily small, but we impose that it vanishes at infinity. We will give a further description below.

Here we note that Einstein equation is written as

$$R_{IJ}^\ell - \frac{1}{2}\bar{g}_{IJ}R^\ell - \Lambda h_{IJ} = (-\bar{g})^{-1/2}T_{IJ},$$

(2.3)

where $h_{IJ} = \bar{g}_{IK}\bar{g}_{JL}h_{KL}$, $R_{IJ}^\ell$ is the linear part of Ricci tensor $R_{IJ}$ with respect to $h_{IJ}$ and all the higher order terms are included in $T_{IJ}$ of the right-hand side. As the left-hand side of eq. (2.3) satisfies the Bianchi identity, we see that

$$\bar{\nabla}_I T^{IJ} = 0$$

(2.4)

holds, where $\bar{\nabla}_I$ is the covariant derivative with respect to $\bar{g}_{IJ}$. Contracting $T^{IJ}$ with the Killing vector of de Sitter space-time,

$$T^I = T^{IJ}\bar{\xi}_J,$$

(2.5)

we obtain the local conservation law

$$\partial_I T^I = 0.$$  

(2.6)

Thus, $\int d^{D-1}x T^0$ is conserved if the surface term $\int d^{D-2}S_i T^i$ vanishes. We define

$$E_{AD} = \frac{1}{8\pi G_D} \int d^{D-1}x T^{0K}\bar{\xi}_K$$

$$= \frac{1}{8\pi G_D} \int d^{D-2}S_i[\bar{\nabla}_I K^{0iJM} - K^{0jJM}\bar{\nabla}_j]\bar{\xi}_J,$$

(2.7)

where $K^{IJKL} = (1/2)(\bar{g}^{IL}H^{KJ} + \bar{g}^{KJ}H^{IL} - \bar{g}^{IK}H^{JL} - \bar{g}^{JL}H^{IK})$, $H^{IJ} = h^{IJ} - (1/2)\bar{g}^{IJ}h_K^K$ and $i = 1, 2, ..., D - 1$. In the case that $\bar{\xi}^I$ is a timelike Killing vector $E_{AD}$ is regarded as the Killing energy, the so called Abbott-Deser(AD) mass.

Next, we rewrite eq. (2.7) in more familiar form in order to obtain its physical meaning. First of all, we need take asymptotic region carefully. In asymptotic flat cases, the extrinsic curvature $K_{ij}$ of slices has the behaviour
\( K^i_j \to 0 \) toward the spatial infinity \( \partial^0 \). As we know, we can expect that such slice does not have asymptotic region in asymptotically de Sitter space-time because static slices with \( K^i_j = 0 \) have no boundary in de Sitter space-time. So we remember that the flat chart of de Sitter space-time has boundary and its spatial metric is conformally flat (See Fig. 2 in ref [10]). Thus, we realize that the most natural condition is given by \( K^i_j \to H \delta^i_j \) or \( K \to (D - 1)H \) [15], where \( H = \sqrt{2\Lambda/(D - 1)(D - 2)} \).

In this chart the metric of the background de Sitter space-time is written as

\[
\bar{g} = -dt^2 + a(t)^2 \sum_{i=1}^{D-1} dx_i^2,
\]

where \( a(t) = e^{Ht} \). The Killing vector has the component \( \xi^M = (-1, Hx^i) \). After simple calculations we obtain

\[
E = a(t) \left[ E_{ADM} + \Delta P_{ADM}(-\xi) \right] = \frac{1}{16\pi G_D} \int d\bar{S}_i(\partial_j h^i_j - \bar{g}^{ij} \partial_j h^k_k) - \frac{1}{8\pi G_D} \int d\bar{S}_i \left[ K^i_j - K \delta^i_j + (D - 2)H \delta^i_j \right] \xi^j,
\]

where \( K^i_j \) is the extrinsic curvature of \( t = \text{constant} \) \( (D - 1) \)-dimensional hypersurface. By using the momentum \( \pi^i_j = K^i_j - q^i_j K, \) where \( q^i_j \) is the metric of \( t = \text{constant} \) hypersurface, the second term in the second line of eq. (2.9) can be written as

\[
\Delta P_{ADM}(-\xi) = P_{ADM}(-\xi) - \bar{P}_{ADM}(-\xi) = -\frac{1}{8\pi G_D} \int d\bar{S}_i (\pi^i_j - \bar{\pi}^i_j) \xi^j.
\]

As a result, \( E = \) is written in terms of the sum of the ADM energy and net momentum. We note that the momentum of background de Sitter space-time is subtracted automatically. The above net momentum is the term ‘momentum’ used in introduction. The above argument does not depends on whether extra dimensions is compactified or not.

Now we compare the above expression with the gravitational energy defined via ‘physical Hamiltonian’ [14]. The gravitational energy has expression

\[
E_{HH} = -\frac{1}{8\pi G_D} \int dS(Nk - \bar{N} \bar{k}) + \frac{1}{8\pi G_D} \int dS(N^i \pi^i_j - \bar{N}^i \bar{\pi}^i_j) r^j,
\]

where \( k \) and \( r^i \) are the trace of the extrinsic curvature and the unit normal vector of \( (D - 2) \)-dimensional surface at the infinity, respectively. \( N \) and \( N^i \) are the lapse function and the shift vector.

The ‘physical Hamiltonian’ is defined by the subtraction the Hamiltonian of a background space-time from the original Hamiltonian. As Hawking & Horowitz showed, the first term of eq. (2.11) is just the ADM energy. Comparing eq. (2.9) with eq. (2.11), one can see that they have the same expression except for scale factor, i.e., \( E = a(t)E_{HH} \), if one chooses the lapse function and the shift vector as follows,

\[
N = \bar{N} = 1 \quad \text{and} \quad N^i = \bar{N}^i = -\xi^i.
\]

It is likely that the priority is given to our definition in asymptotically de Sitter space-times because the Hawking & Horowitz’s construction of physical Hamiltonian has the operation of artificial subtraction to keep it finite.

In asymptotically flat cases, the ADM energy is defined by view point of a static observer with \( N = 1 \) and \( N^i = 0 \) at spatial infinity \( \partial^0 \). The naturalness of observer selection to define the energy is related to the timelike translation symmetry at infinity. On the other hand, in asymptotically de Sitter space-time, we must consider hypersurfaces which have infinity in order to define the non-zero energy. As we have done, the most convenient slices is one which corresponds to the flat slices in de Sitter space-time. This slices reach into the timelike infinity \( I^+ \) (See Fig. 2 in Ref. [10]). However, the slices is not associated with the timelike translation symmetry which de Sitter space-time possesses, that is, the slice is not orthogonal to timelike Killing vector. Thus, it may be natural that the ADM momentum term enters into the expression of the mass.

In four dimension the AD mass can be negative when system has a large outward momentum [10]. On the other hand, the positivity of the AD or ADM mass is guaranteed for systems without the net momentum, \( \Delta P_{ADM} = 0 \), in four dimension [16].

III. EXAMPLES IN FIVE DIMENSIONS

In the same way as four dimensional cases, the AD mass can be negative regardless of the compactification of the extra dimension if the system has a large momentum. However, the physical reason of the negativity is the dynamics
of four dimensional part of space-time, rather than the (quantum) stability. Now we are interest in just stability. So, we consider the situation only where the contribution of the net ADM momentum does not exist. We give two examples in five dimension which are regular everywhere and have negative energy.

Let us consider an initial slice with \( K_{ij} = \pm H g_{ij} \) and \( H = \sqrt{\Lambda/6} \). In this slice the Hamiltonian constraint becomes \( ^{(4)}R = 0 \). Thus one can use the argument on the momentarily static slices in asymptotically flat cases because the Hamiltonian constraint is just same one.

One can see easily that the euclidian Reissner-Nordstrom metric with imaginary ‘charge’ i.e satisfies the Hamiltonian constraint \( ^{(4)}H \). This metric of the hypersurface is given by

\[
(4)g = V(r)d\chi^2 + \frac{dr^2}{V(r)} + r^2d\Omega_2^2
\]

where \( V(r) = 1 - 2m/r - e^2/r^2 \) and \( r \geq r_+ := m + \sqrt{m^2 + e^2} \). To avoid a conical singularity at \( r = r_+ \), we assume the period \( \chi_0 = 4\pi/V'(r_+) = 2\pi r_+^2/(r_+ - m) \) along the \( \chi \)-direction.

On the present slice, the mass is constructed by only the ADM energy component, \( E_{AD} = E_{ADM} = m\chi_0/2G_5 = m/2 \). By the same argument as Brill & Horowitz \( ^{(4)}H \), it is shown that the mass becomes negative and does not have the lower bound. The mass can be set to be arbitrary negative regardless of the radius of the compactified space. However, it is not obvious whether the AD energy is conserved or not in the course of time development because of the existence of the cosmological constant and momentarily non-static slices. In fact we will see that the energy is not conserved in the next example.

As a second example, we take a dynamical solution \( ^{(5)}H \). The metric is

\[
ds^2 = -dt^2 + a(t)^2 \left[ \frac{d\chi^2}{\Delta} + \Delta dr^2 + r^2\Delta^2 d\Omega_2^2 \right],
\]

where \( \Delta(r) = 1 - m/r, a(t) = e^{\pm Ht} \) and \( H = \sqrt{\Lambda/6} \). The \( t \)-constant hypersurface has the extrinsic curvature \( K^t_j = \pm H\delta^t_j \). However, this space-time has timelike naked singularity at \( r = m \). Since the expansion of outgoing null geodesics congruence is

\[
\theta_\pm = \frac{1}{\sqrt{2}} \left( \pm 2H + \frac{4r - m}{2a\Delta^{3/2}r^2} \right),
\]

the apparent horizon does not exist in the expanding chart with \( a = e^{Ht} \). In collapsing chart with \( a = e^{-Ht} \), apparent horizon also does not exist although surfaces such that \( \theta_\pm = 0 \) exists. The curvature invariant is given by

\[
R^{tJKL}R^{tJKL} = 5H^4 + 3 \left[ H^2 - \frac{m^2}{2a^2(r-m)^3} \right]^2 + \left[ H^2 - \frac{m}{a^2(r-m)^3} \right]^2.
\]

To avoid the naked singularity at \( r = m \), we change the sign of the mass parameter, \( m \rightarrow -m \). After that the radial coordinate \( r \) can run up to \( r = 0 \) and the conical singularity occurs at \( r = 0 \) in general cases. Near \( r = 0 \), the metric is written as

\[
ds^2 \simeq -dt^2 + a(t)^2 \left[ \frac{r}{m}d\chi^2 + \frac{m}{r}dr^2 + m^2d\Omega_2^2 \right].
\]

Here we introduce a new coordinate \( R = (rm)^{1/2} \) and the metric is

\[
ds^2 \simeq -dt^2 + a^2 \left[ 4\left\{ dR^2 + R^2d\left( \frac{\chi}{2m} \right)^2 \right\} + m^2d\Omega_2^2 \right].
\]

Hence, the metric is regular everywhere except for a \( a = 0 \) singularity if one assumes the period \( \chi_0 = 4\pi m \) along the \( \chi \)-direction. The physical size of the compactified dimensions is given by \( 4\pi a(t)m \). Thus, it decreases/increases if one takes the collapsing/expanding chart. A cosmological constant make the compactified space dynamical as well as the four dimensional part. The AD mass is

\[1\]Here we used the fact that the five dimensional Newton’s constant is written as \( G_5 = \chi_0 G_4 \) in term of the four dimensional one.
\[ E = E_{ADM} = -\frac{a(t)^3 m \chi_p}{G_5} = -a^3(t) m. \] (3.7)

This is not conserved due to the expansion of the universe! This comes from the non-vanishing boundary term \( \int d^{D-2} S_i T^i \) which vanishes for asymptotically flat cases. At first glance, best we can do might be choosing the collapsing chart and we can keep the radius of the compactified space to be less than Planck length. In the chart, we observe the value of the energy approaches zero. One may think that a large cosmological constant stabilises collapsing chart and we can keep the radius of the compactified space to be less than Planck length. In the chart, not covariant for full space-times. The excellent approach bearing supergravity in mind was given in ref. [19]. However, both with the large compactified space is suppressed.

The vector \( \xi' = \bar{\epsilon}' \gamma_5 \epsilon \) defined by spinor \( \epsilon \) satisfying \( \nabla_i \xi = 0 \) is conformal Killing vector, not Killing vector in de Sitter space-time. That it, \( \bar{\epsilon}' \) is a conformal Killing spinor, not Killing spinor. This point is the main reason why one cannot prove the positivity of the AD mass for asymptotically de Sitter space-times having the net momentum because the left-hand side of eq. (4.4) cannot be the AD mass.

IV. POSITIVE ENERGY THEOREM, ELEMENTARY FERMION AND STABILITY

In locally asymptotically flat cases with compactified dimensions, the break down of positive energy theorem means non-existence of the Witten spinor [2]. In supergravity side, that means there are not supersymmetry because the spinor is related to infinitesimal generator of local supersymmetry [18].

In asymptotically de Sitter cases, the situation is different from the above. As we stated, the energy can be negative even if the extra dimension is not compactified and the topology is trivial. In this section, we discuss the positive energy theorem, based on [19] [19] in cases whose extra dimension is not compactified. Conversely, we can see easily that the Witten spinor cannot exist for examples given in the previous section.

Following Kastor & Traschen [19], we define the cosmological supercovariant derivative operator on a spinor \( \gamma \) as

\[ \hat{\nabla}_I \gamma = \left( \nabla_I + \frac{1}{2} H \gamma_I \right) \gamma = \left( \partial_I + \Gamma_I + \frac{1}{2} H \gamma_I \right) \gamma =: \left( \partial_I + \Gamma'_I \right) \gamma, \] (4.1)

where \( \Gamma_I \) is the spin connection,

\[ \Gamma_I = -\frac{1}{8} \epsilon^J K \nabla_J \epsilon^I [\gamma_L, \gamma_K], \] (4.2)

and \( \epsilon^I \) is quasi-orthogonal basis. The cosmological Witten equation [1] is define by

\[ \gamma^I \hat{\nabla}_I \gamma = 0. \] (4.3)

The solution is given by a constant spinor \( \epsilon_0 \) satisfying \( \gamma^0 \epsilon_0 = -i \epsilon_0 \) in the expanding flat slice of de Sitter space-time.

By using the Bianchi identity, the Hamiltonian and momentum constraints, we obtain an identity

\[ \int dS^i \epsilon^I \hat{\nabla}_I \epsilon = \int dV [\left| \hat{\nabla}_I \gamma \right|^2 + 4 \pi G_D (\epsilon^I \epsilon_0 + \epsilon^I T^0_0 + 1^I \gamma^I \gamma^0 \epsilon)] \] (4.4)

Let us evaluate the left-hand side. We follow Witten’s argument [21] carefully. First, the constant spinor satisfies

\[ \gamma^I \hat{\nabla}_I \epsilon_0 = \gamma^I \Gamma'_I \epsilon_0 = \gamma^i(D-1) \Gamma_i \epsilon_0 + \frac{i}{2} [-K + (D - 1) H] \epsilon_0 \] (4.5)

The argument in ref. [19] is prototype of the proof. Rigourously speaking, the ‘supercovariant derivative’ defined there is not covariant for full space-times. The excellent approach bearing supergravity in mind was given in ref. [19]. However, both papers have not been attained in the present refined form. The explicit evaluation of the left-hand side of eq. (4.4) has not done yet in ref. [19]. See ref. [20] for asymptotically AdS space-times.
where \((D-1)\Gamma_i = -(1/8)\epsilon j^{k(D-1)}\nabla_i e^k[j_i, j_k]\) and we used
\[
\Gamma'_i = (D-1)\Gamma_i + \frac{1}{2}K_{ij}\gamma^i\gamma^0 + \frac{i}{2}H\gamma_i \tag{4.6}
\]
If we suppose \(K = (D - 1)H + O(1/r^{D-1})\), the same argument of the existence of solution as asymptotic flat holds and we obtain
\[
\int dS'\epsilon^i \nabla_i \epsilon = \int dS'\epsilon^i_0(\Gamma'_i - \gamma_i\gamma^0\Gamma'_0)\epsilon_0. \tag{4.7}
\]
Inserting the decomposition \([4.4]\) of the spin connection into the above, we obtain the familiar expression
\[
\int dS'\epsilon^i \nabla_i \epsilon = \int dS'\epsilon^i_0((D-1)\Gamma_i - \gamma_i\gamma^{(D-1)}\Gamma_j)\epsilon_0 + \frac{1}{2} \int dS'\epsilon^i_0[K_j^i - \delta_j^iK + (D - 2)H\delta_j^i] \gamma^j\gamma^0 \epsilon_0
\]
\[
= \frac{1}{4} \int dS'\epsilon^i_0(\partial_j h^j_i - \partial_i h^j_j)\epsilon_0 - \frac{1}{2} \int dS'_iK^i_j\epsilon_0 \gamma^j \epsilon_0
\]
\[
= 4\pi G_D \left(E_{ADM}\epsilon_0^2 + \Delta P_{ADM}(\epsilon_0^2\epsilon_0)\right), \tag{4.8}
\]
where \(K^i_j\) is the traceless part of \(K^i_j\). Here we note that \(\epsilon_0^i\gamma^i \epsilon_0 = -e_0^i|\epsilon_0|^2 = O(1/r^{D-3})\) due to \(\gamma^0 \epsilon_0 = -\epsilon_0\). This and \(K^i_j = O(1/r^{D-2})\) lead us that the net momentum term \(\Delta P_{ADM}(\epsilon_0^2\epsilon_0)\) vanishes. Finally we obtain an inequality
\[
E = \frac{a(t)}{4\pi G_D} \int dS'\epsilon^i \nabla_i \epsilon
\]
\[
= a(t)E_{ADM}|\epsilon_0|^2 \geq 0 \tag{4.9}
\]
under the dominant energy condition \([23]\) on the energy-momentum tensor \(T_{IJ}\). Unfortunately, we cannot say something about \(E = 0\) case because Witten spinor approaching to \(\epsilon_0\) at infinity is uniquely determined\([14]\). Thus, there is not contradiction with the existence of non-trivial solution which has zero ADM mass and is given in ref. \([14]\).

We can see that \(E\) is not equal to the AD mass, \(E_{AD}\). When \(K^i_j = O(1/r^{D-1})\) holds and the momentum term of the AD mass vanishes, the AD mass equals to the ADM energy and the positivity is guaranteed. For our purpose, it is worth imposing \(K^i_j = O(1/r^{D-1})\) because we are interest in the stability of space-time, not its dynamics. As eq. \([4.9]\) this implies the positivity of the AD mass. The apparent contradiction with examples given in the previous section indicates that the Witten spinor does not exist in such examples.

V. SUMMARY

In this paper we defined the gravitational mass in asymptotically de Sitter space-time with extra dimensions and obtained a refined expression related to the ADM energy and momentum associated to the timelike Killing vector of the background de Sitter space-time. Furthermore, we gave one dynamical solution and one initial data with the negative energy in five dimensions. Does these solution indicate the quantum decay of de Sitter space-time? We cannot reply to the question instantly because we do not know whether the instanton exists or not. Naively speaking, we can guess from the previous section that the decay occurs unless one imposes the existence of elementary fermion or supersymmetry.

Finally, we should stress the fact that the energy seems to depend on the time in general although the definition is reasonable. This means that the contribution from the boundary is not negligible.

**In asymptotically flat space-times, \(\epsilon_0\) is arbitrary constant spinor. This means the existence of D's independent solution of Witten equation. As a result, we obtain \(R_{ijkl} = 0\) from 0 = \([\nabla_i, \nabla_j]\epsilon = (1/4)R_{ijkl} [\gamma^k, \gamma^l] \epsilon.\)
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