Global monopole asymptotic solutions in Hořava gravity

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Abstract
In Hořava’s theory of gravity coupled to a global monopole source, we seek for static, spherically symmetric spacetime, solutions for general values of \( \lambda \). We obtain the explicit solutions with deficit solid angles, in the IR modified Hořava gravity model, at the IR fixed point \( \lambda = 1 \) and at the conformal point \( \lambda = 1/3 \). For the other values of \( 1 > \lambda > 0 \) we also find special solutions to the inhomogenous equation of the gravity model with detailed balance, and we discuss a possibility of astrophysical applications of the \( \lambda = 1/2 \) solution that has a deficit angle for a finite range.

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Recently, Hořava proposed a renormalizable gravity theory in the UV limit [1], a lot of related works have been widely circulated. Studies on Hořava–Lifshitz cosmology [2], black hole solutions [3, 4] and other interesting topics [5] have been reported. It was also shown in [6] that classical solutions to the IR limit of Hořava gravity can mimic general relativity plus cold dark matter. On the other hand, authors of [7] suggested that static vacuum solutions in projectable cases of Hořava gravity (in the IR limit) would not be capable of explaining galactic dark matter problem. We consider global monopole (GM) as possible dark matter in this gravity with relaxing the projectability condition.

In the IR modified Hořava theory of gravity [1] where the detailed balance condition is softly violated (via the term proportional to \( \omega R \) in equation (4) below), we study geometric structures affected by gravitationally coupled GM source [8]. Considering static, spherically symmetric spacetimes, we obtain solutions to a set of equations derived for general values of \( \lambda \). In this IR modified model, we find the explicit solutions at the IR fixed point \( \lambda = 1 \) and at the conformal point \( \lambda = 1/3 \). In both cases deficit solid angles occur.

For the other values of \( 1 > \lambda > 0 \) in the case with detailed balance \( \omega = 0 \), we have new special solutions, in addition to known general solutions [3] to the corresponding homogeneous equation. By simple analysis, we show that the GM spacetime in the case \( \lambda = 1/2 \) can have...
a deficit solid angle only for a finite range and that it is asymptotically flat. We discuss a possibility of its astrophysical applications.

Using the ADM decomposition of the spacetime metric
\[ ds^2 = -N^2 dt^2 + g_{ij}(dx^i + N^i dr)(dx^j + N^j dr) \]  
(1)
with the lapse \( N \) and shift fields \( N^i \), the IR modified Hořava gravity theory is described by the action
\[ S_H = \int dt d^3 x \sqrt{g} N (L_K + L_V), \]  
(2)
where the kinetic term
\[ L_K = 2\kappa^{-2} (K_{ij}K^{ij} - \lambda K^2) \]  
(3)
is made of the extrinsic curvature \( K_{ij}(\equiv 2N^{-1}(\nabla_i N_j - \nabla_j N_i)) \), its trace \( K \) and a parameter \( \lambda \). \( L_V \) includes all potential terms satisfying the detailed balance condition \([1]\), and is given by
\[ L_V = \frac{\mu^2}{2(1 - 3\lambda)} R^2 + \Lambda R - 3\Lambda^2 - \frac{\mu^2 \omega}{8(1 - 3\lambda)} R, \]  
(4)
where \( \Lambda (\lessgtr 0) \) is a cosmological constant. The last term (\( \propto \omega R \)) which violates softly the detailed balance condition is added so that it can lead to a Minkowski vacuum (when taking \( \Lambda \to 0 \)) and have Schwarzschild analog solutions in the IR (see equation (15)) \([1, 9]\].

Following \([3]\), we relax the projectability condition (with \( N(t, \vec{x}) \) in equation (1)) and take a broader view of Hořava \([1]\) in which the relative coefficients of the ADM decomposition of the Einstein–Hilbert action are modified and additional terms of a spatial higher derivative are included in equations (2)–(4). To seek for static, spherically symmetric spacetime solutions, we adopt the metric ansatz as
\[ ds^2 = -N^2(r) dt^2 + dr^2 + r^2(d\theta^2 + \sin^2\theta\,d\phi^2), \]  
(5)
and we can write \( S_H = 4\pi \int dt dr \mathcal{L}_H \) with the Lagrangian density
\[ \mathcal{L}_H = \frac{N}{q^2 \sqrt{f}} \left[ 3\Lambda^2 r^2 + 2(\omega - \Lambda)(1 - f - rf') \right. \]  
\[ + \left. \frac{(1 - \lambda)f'^2}{2} + (1 - 2\lambda)(1 - f)^2 r^2 - 2\lambda(1 - f)f' r \right], \]  
(6)
where \( q^2 = 8(3\lambda - 1)/(\kappa^2 \mu^2) \).

Let us consider a GM source which has the action written up to second order in space derivatives \( \mathcal{O}(\partial_j \Phi^2) \) \([10]\):
\[ S_{\text{matter}} = -\int dr d^3 x \sqrt{g} N \left[ -\frac{1}{2N^2} \partial_i \Phi \cdot \partial_i \Phi + \frac{1}{2} g^{ij} \partial_i \Phi \cdot \partial_j \Phi + \frac{\chi}{4} (\Phi^2 - \eta^2)^2 \right], \]  
(7)
with a dimensionless coupling constant \( \chi \). We can write \( S_{\text{matter}} = 4\pi \int dt dr \mathcal{L}_{\text{matter}} \) with
\[ \mathcal{L}_{\text{matter}} = -\frac{N r^2}{\sqrt{f}} \left[ \frac{1}{2} \left( f h^2 + \frac{2h^2}{r^2} \right) + \frac{\chi}{4} (h^2 - \eta^2)^2 \right], \]  
(8)
where a hedgehog ansatz for the GM, \( \Phi = h(r)\vec{x}/r \), is assumed.
Performing a variation of the total action $S_{\text{total}} = S_H + S_{\text{matter}}$ with respect to $h(r)$, $N(r)$, and $f(r)$ respectively, we obtain the following equations:

$$\frac{\sqrt{f}}{N r^2} \left( \frac{N r^2}{\sqrt{f}} f' \right)' = \frac{2h}{r^2} + \chi (h^2 - \eta^2)h,$$

$$\left(1 - \lambda\right) \frac{f''}{2} + (1 - 2\lambda) \frac{(1 - f)}{r^2} = -2\lambda \frac{f'}{r} + 2(\omega - \Lambda)(1 - f') + 3\Lambda^2 r^2$$

$$= q^2 r^2 \left[ \frac{h'^2}{2} + \frac{h^2}{r^2} + \frac{\chi}{4} (h^2 - \eta^2)^2 \right].$$

$$\left( \frac{N}{\sqrt{f}} \right)' \left[ -2\lambda \frac{(1 - f)}{r} + (1 - \lambda) f' - 2r (\omega - \Lambda) \right]$$

$$= - \frac{N}{\sqrt{f}} \left[ 2(1 - \lambda) \frac{(1 - f)}{r^2} + (1 - \lambda) f'' + q^2 r^2 h^2 \right].$$

It was suggested in [11] that the black hole no-hair conjecture [12] can be violated for GM (which is one kind of topological defect), and the GM solution to equation (9) derived from $S_{\text{matter}}$ including the Mexican hat potential in equation (7) is (when neglecting $O(1/r^2)$ [13, 14])

$$h(r) = \eta,$$

which is valid for the outside of the GM core, $r > \chi^{-1/2} \eta^{-1}$ [8]. With the asymptotic solution (12) for $h(r)$, the solutions to the other equations (10) and (11) for various values of $\lambda$ are given as follows.

1. $\lambda = 1$ case

In this $\lambda = 1$ case where Ho\u017ava’s theory coincides with Einstein’s general theory of relativity in the IR limit, equation (12) gives us the simple solution to equation (11) as $N/\sqrt{f} = 1$. Putting $1 - f \equiv -(\omega - \Lambda) r^2 + X^{1/2}$, we can rewrite the remaining equation (10) as

$$3\omega (\omega - 2\Lambda) r^2 + q^2 \eta^2 = \frac{X'}{r} - \frac{X}{r^2},$$

whose solution is

$$f = 1 + (\omega - \Lambda) r^2 - \sqrt{\omega (\omega - 2\Lambda) r^4 + q^2 \eta^2 r^2 + \beta r}.$$  

Note that equation (14) would be the same as the result of [15] if there were not the additional new term $q^2 \eta^2 r^2$. In the limit $r \gg \sqrt{q^2 \eta^2 / [\omega (\omega - 2\Lambda)]}$ and $r \gg [\beta / [\omega (\omega - 2\Lambda)]]^{1/3}$, equation (14) can be approximated as

$$f = 1 - \frac{q^2 \eta^2}{2 \sqrt{\omega (\omega - 2\Lambda)}} + \frac{\Lambda_{\text{eff}}}{2} r^2 - \frac{\beta}{2 \sqrt{\omega (\omega - 2\Lambda)} r},$$

which can be compared with the Schwarzschild-AdS black hole carrying a GM charge $q \eta$ and a mass $M \simeq \beta / [4 \sqrt{\omega (\omega - 2\Lambda)}]$. Here an effective cosmological constant $\Lambda_{\text{eff}} \equiv 2[(\omega - \Lambda) - \sqrt{\omega (\omega - 2\Lambda)}](\simeq \Lambda^2 / \omega$ for $-\Lambda < \omega$) and a deficit angle $q^2 \eta^2 / [2 \sqrt{\omega (\omega - 2\Lambda)}] = 8\eta^2 / [\kappa^2 \mu^2 \sqrt{\omega (\omega - 2\Lambda)}].$
2. $1 > \lambda$ case

With equation (12) in the case where $1 > \lambda$, we can rewrite equation (10) as

\[
\frac{1 - \lambda}{2} \left( \frac{dY}{du} \right)^2 + \frac{1 - 3\lambda}{1 - \lambda} Y^2 = 3\omega(\omega - 2\Lambda)e^{2\gamma(1 + \mu)u} + q^2\eta^2 e^{(1 + \mu)u},
\]

with $u = \ln r$ and $Y(u(r)) = r^{2/3}/(1 - \lambda)$.

2.1. $\lambda = \frac{1}{3}$ case

In the case $\lambda = \frac{1}{3}$ where it is allowed for us to get a nontrivial conformal limit [16], the set of equations in equations (9)–(11) obtained from equations (4) and (8) can be replaced by the same form with only substitution $q^2 \rightarrow q^2/(3\lambda - 1) = 8/(\kappa^2\mu^2)$, and we have their solutions

\[
f = 1 + (\omega - \Lambda)r^2 - \frac{2M}{r} - \frac{\sqrt{3}}{9\omega(\omega - 2\Lambda)r} \left[ \frac{8\eta^2}{\kappa^2\mu^2} + 3\omega(\omega - 2\Lambda)r^2 \right]^2,
\]

but $N(r)$ is not given by a simple expression. Equation (17) goes to

\[
f = 1 - \frac{4\eta^2}{\kappa^2\mu^2\sqrt{\omega(\omega - 2\Lambda)}} + \frac{\lambda_{\text{eff}}}{2}r^2 - \frac{2M}{r},
\]

in the region $r \gg \sqrt{8\eta}/[\kappa\mu\sqrt{\omega(\omega - 2\Lambda)}]$. In this large $r$ limit, $f(r)$ (of this case $\lambda = 1/3$) is almost the same as one of the $\lambda = 1$ case (in equation (15)) except different values of the deficit angle, while the lapse functions $N(r)$ in these cases are very different from each other.

When there is no GM source, we have

\[
f = 1 + \lambda_{\text{eff}}r^2 - \frac{2M}{r},
\]

and $N^2 = f(r)$.

2.2. $1 > \lambda > \frac{1}{3}$ case

From now on, we consider the case ($1 > \lambda > \frac{1}{3}$) with detailed balance condition (i.e. $\omega = 0$) for simplicity and without $q^2$ rescaling which is done in the case $\lambda = 1/3$.

Equation (16) can be written as a simple inhomogeneous equation

\[
\left( \frac{dY}{dU} \right)^2 = AY^2 + Bu,
\]

where $U \equiv \gamma u$, $A = 2(3\lambda - 1)/(\gamma^2(1 - \lambda)^2)$, $B = 2q^2\eta^2/(\gamma^2(1 - \lambda))$, and $\gamma = 2(1 + \lambda)/(1 - \lambda)$.

The solution is

\[
f = 1 - \lambda r^2 - q\eta\sqrt{\frac{1 - \lambda}{3\lambda - 1} \frac{r}{R(r)}},
\]

where with a constant $c_R$,

\[
\ln r(R) \left( \frac{U}{\gamma} \right) = c_R - \sqrt{2(3\lambda - 1)} \ln \frac{1 + R^2 - 1}{R} + \frac{1 + \lambda}{\lambda - 3} \ln \left[ 1 - \frac{\sqrt{2(3\lambda - 1)(1 + R^2)}}{1 + \lambda} \right].
\]

1 This range is known to admit a sort of ghost-like state [17], but with relaxing the projectability condition there seems to be no problem as suggested in [18].
From the last equation we can estimate the asymptotic behavior of equation (21) as $r/R \approx 0$ for large $R$, while, for small $R$, $r/R \propto r^{1-1/n(\lambda)}$ with $n(\lambda) \equiv -1 + (4 + \sqrt{2(3\lambda - 1)})/(3 - \lambda)$ and $-1 < 1 - 1/n(\lambda) < 1/2$. Since especially

$$\frac{r}{R(r)} = \delta(\text{constant})$$

(23)

for $R^2 < 5/4$ when $n(\lambda = 1/2) = 1$, we may have a deficit angle for a finite range $r < r_0(\equiv \sqrt{5\delta}/2)$ in the case $\lambda = 1/2$.

The lapse function $N(r) = \sqrt{f(r)L(r)}$ with

$$L(r) \equiv \frac{2\lambda}{(1-\lambda)r} - \frac{2(3\lambda - 1)(1 - f - \Lambda r^2)}{(1-\lambda)r^2\sqrt{2(1-\lambda)q^2\eta^2 + 2(3\lambda - 1)(1 - f - \Lambda r^2)^2/r^2}}$$

(24)

When $q \eta = 0$, the lapse [3] which is a solution to the corresponding homogeneous equation of (20) (instead of the last term in equation (21)).

2.3. $\lambda < 1/3$ case

When $\lambda < 1/3$, equation (20) is replaced by

$$\left(\frac{dY}{dU}\right)^2 = -\alpha Y^2 + Be^U,$$

(26)

with $\alpha = 2(1 - 3\lambda)/(y^2(1 - \lambda)^2) > 0$ and $B > 0$ given below equation (20). This inhomogeneous equation has a (special) solution

$$f = 1 - \Lambda r^2 - q\eta \sqrt{\frac{1-\lambda}{1-3\lambda}} I(r),$$

(27)

where

$$\ln r(I) = \frac{U}{\gamma} = c_I - \frac{\sqrt{2(1-3\lambda)}}{\lambda - 3}\arctan \frac{I}{\sqrt{1 - I^2}} + \frac{1+\lambda}{\lambda - 3} \ln |\frac{2(1-3\lambda)(1 - I^2)}{1 + \lambda}|$$

(28)

The lapse function in this case can be obtained by the similar method as we have done in equation (24).

In summary, we have studied the IR-modified Hořava theory of gravity. In static, spherically symmetric spacetimes, we obtain asymptotic solutions (valid outside a GM core) for general values of $\lambda$ to the equations of gravity coupled to the GM. As we can see from equations (15) and (18) obtained in the large $r$ limit, in the cases $\lambda = 1$ and $\lambda = 1/3$ we have deficit angles as in Einstein’s theory of gravity coupled to the GM [8]. $f(r)$ of the case $\lambda = 1/3$ (in equation (18)) is almost the same as one of the $\lambda = 1$ case (in equation (15)) except different values of the deficit angle. We also have the explicit solution of the lapse function in the $\lambda = 1$ case.

In the case $1 > \lambda > 1/3$ we have studied the Hořava model with detailed balance and obtained special solutions, in addition to known general solutions [3] to its homogeneous
equation. When especially $\lambda = 1/2$, $f = 1 - q\eta\delta - \Lambda r^2$ for $r < r_0(= \sqrt{3}\delta/2)$ as seen in equations (21) and (23), and we can have a GM spacetime that has a deficit solid angle for a finite range (and is asymptotically flat if $\Lambda \to 0$), which is different from the GM spacetime in Einstein’s theory of gravity [8]. This might be more helpful for us, with the GM as [14, 19], to explain near flatness of rotation curves of galaxies, which appears over a finite range $0 \ll r < r_0$ (when we take the constant $r_0$ at about ten times larger value than a typical galaxy size).

To explain the near flatness of rotation curves in preceding models using GM with an energy density proportional to $r^{-2}$, we need nonlinear coupling between gravity and the GM as nonminimal coupling in [19] or Brans–Dicke field coupling. In the latter case, as discussed below equation (4) of [20], it can be yielded by the finite range, logarithmic gravitational potential that is derived from the Brans–Dicke field equation. For the rotation velocity formula to be valid only for the finite range given by the (galactic halo) radius $r_0$, the responsible GM field should vanish at distance larger than $r_0$ due to interactions with the nearest topological defect such as anti-monopole, in the way that the GM field lines can be absorbed into the anti-monopole core, as argued in [20].

Actually, from equations (21)–(25) and the rotation formula $v_{rot} = \sqrt{r(fL^2)/(2fL^2)}$ [14] we have a contribution $\sqrt{2q\eta\delta/5}R$ (for $R < \sqrt{5}/2$ and $q\eta\delta \ll 1$) which might be compared with a measured value of stellar rotation velocity $100 - 300\,\text{km/s}$, even if we cannot prove because of the nonlinear property of equation (20) that all other contributions from solutions to inhomogeneous and homogeneous equations (which are given in equation (25) and below it with $\lambda = 1/2$) than the contribution cancel out. By further studies of nonlinear equation (20) or by considering Brans–Dicke field coupling to Hořava gravity [21] as preceding models [14, 19], with the $\lambda = 1/2$ Hořava gravity solution given in equation (23) having a finite range deficit angle, we anticipate that more natural explanations for the near flatness can be possible. This kind of $r$-dependent, deficit solid angle was obtained in Brans–Dicke gravity theory [22], by studying the quantum effects [23] due to the GM, which can be expressed as quadratic in curvature as if the Hořava gravity with detailed balance.

We add a comment for the case $\lambda < 1/3$. A small $\tilde{F} \ll (1 - 3\lambda)/(3(1 - \lambda)(1 - \lambda/3))$ approximation for $r(I)$ in equation (28), $I \propto 1 + \ln(r/r_0)$, gives us a $r$-linear contribution from equation (27) to the lapse $N^2 = fL^2$. The contribution induced by GM might be related with such a ($r$-linear) gravitational potential $\psi (= (N^2 - 1)/2)$ as suggested by the author of [24] for the purpose of explaining stellar rotation curves, which show a tendency to rise at large distances. (According to the reference, closer examination of the rotational velocity data reveals that some of rotation curves are growing universally at large distances as $v_{rot}^2 \propto r$.) This tendency can arise from our $r$-linear gravitational potential as well as from the potential of [24], as can be easily seen by applying the formula $v_{rot}^2 = \tilde{r} \cdot \tilde{\nabla}\psi$ [14]. We would like to make just one more comment on GM special solutions with $\lambda < 1$: the seeds from which galactic dark matter and large-scale structure were formed in the early universe could be generated at GUT UV scale in topological defect models [25], and their remnants could remain as parts of the present galactic halo responsible for the nearly flatness (or even the rising tendency) of outer rotation curves. We thus have studied special solutions not only in the case $\lambda = 1$ but also $\lambda < 1$. The latter case might not follow the requirement of invariance under all spacetime diffeomorphisms but improve the UV behavior [1] while playing the possible role of galactic dark matter. When we almost complete our study, we see [26] that has reported results including some information consistent with ours in section 1. We have not considered in equation (7) higher derivative terms of GM fields in Hořava–Lifshitz theory for simplicity. Even if we add these terms $(\partial_j(\partial_k\Phi^e)^{(2-1)/2}\tilde{\Phi})^2 (1 < z < 3)$ [10, 26] which is $O(1/r^{4+2c})$, our
main results are not changed in the large $r$ limit while neglecting terms of the same order as ones neglected when we have got the vacuum solution in equation (12).

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