Introduction. — General relativity (GR) has been a successful theory of gravity for explaining the motion of planets and stars at the macroscopic scale so far since it is strongly valid under the weak gravitational-field approximation. However, the theory has faced on theoretical incompleteness in terms of unified theory with quantum nature, for example, describing physics at strong gravitational-field near astrophysical compact objects such as black holes or neutron stars, mysterious darkness such as dark energy/matter, the origin of the Universe such as inflationary blow-up and hierarchical phase transition at the very early stage of the Universe (see §1 and references therein). Despite many trials for introducing new particles and/or alternative gravity theories have been suggested in order to address these issues, a remarkable breakthrough is still required in the theoretical as well as experimental/observational point of view.

One of discrepancies between gauge theories and GR is that GR is not renormalizable due to ultraviolet (UV) divergence. The UV divergence appears in the extreme circumstances and/or at high energies, where the quantum effect should be taken into account, for example, at the early stage of the Universe or in the vicinity of a black hole. A renormalizable gravity theory is naturally supposed to have improved properties at the UV scale. One of theories containing such philosophy has been proposed by Hořava [8,10] to make a UV complete theory through sacrificing local Lorentz symmetry, which is called Hořava-Lifshitz (HL) gravity. The HL gravity approximates to GR at infrared (IR) scale while it becomes a different type of gravity at UV scale by introducing an anisotropic scaling between time and space. This anisotropic scaling breaks the local Lorentz symmetry but it is the key of the renormalizability in HL gravity. There have been many intensive studies in diverse fields of applications as a possible candidate of quantum gravity (for recent progress reports, see Refs. [11–13] and references therein).

Meanwhile, the deformed HL gravity has been introduced to get asymptotically flat solutions by introducing a parameter ω in addition to parameters in the original HL gravity [9]. A spherically symmetric and asymptotically flat solution has been found by Kehagias and Sfetsos (KS) [14], approaching Schwarzschild black hole (SBH) at IR limit. Among several parameters in the deformed HL gravity, the only free parameter is ω in the KS solution. The KS solution in the deformed HL gravity has been constrained by several observations in Refs. [15,17], by introducing a dimensionless parameter \( \tilde{\omega} = \omega(GNc^{-2}M)^{2} \). The light deflections by the Sun, the Jupiter and the Earth give the low bounds \( \omega_{\min} \sim 10^{-15} \), \( 10^{-17} \) and \( 10^{-16} \), respectively [15]. The range-residual of the Mercury in the solar system and the orbital motion of the S2 star [18] around the supermassive black hole of mass \( 4 \times 10^{6}M_{\odot} \) at the center of the Milky Way constrain HL gravity by \( \omega \sim 10^{-8} \) [16]. The weak lensing by galaxies of mass \( 10^{10}M_{\odot} \) reads the low bound \( \omega_{\min} \sim 10^{-16} \) [17]. These low bounds are converted to \( \omega_{\min} \sim 10^{-48} \) -- \( 10^{-16} \) cm\(^{-2} \).

Here one might ask a question of how extreme circumstances can affect the structural equilibrium between HL gravity and isotropic matters. The same question in GR has been arisen by Tolman [19] and Oppenheimer and Volkoff [20] (TOV). They formulated so-called the TOV equation to describe the equilibrium state of compact stars like white dwarf and neutron star (NS) in spherically symmetric and static configuration. By solving the TOV equation, we can obtain the mass-radius relation and, therefore, can estimate the maximum mass of a compact star and its radius. In particular, for NS, the maximum mass and radius are sensitive to the selected equation-of-state (EoS) model.

In this Letter, we investigate the structure of NS in
the HL gravity. The TOV equation of perfect fluids in HL gravity is derived first for the purpose. Then, in turn, we solve the equation with several EoS models of Akmal-Pandharipande-Ravenhall (APR) [21], Mütter-Prakash-Ainsworth (MPA1) [22], Müller-Serot (MS1) [23], and Wiringa-Fiks-Fabrocini (WFF1) [24] since, from TOV calculation in GR, these models have shown consistent results to the observed maximum mass, \( \sim 2M_\odot \), where \( M_\odot \) represents the solar mass [25][26].

For the numerical computation of TOV equation, we adopt the fifth order Runge-Kutta solver [27][28]. The mass-radius relation in HL gravity is formulated by an anisotropic scaling shift vector, and the topologically massive gravity action \([29, 30]\). The HL gravity is formulated by an anisotropic scaling shift vector, and the topologically massive gravity action \([29, 30]\). The TOV equation of perfect fluids in HL gravity is derived first for the purpose. Then, for the mass-radius relation in HL gravity as well as the parametric limit of a HL parameter, \( \omega \), in terms of TOV analysis. Finally, we discuss our results in terms of observational validation.

**TOV Equation in HL gravity.**

Let us begin with the Arnowitt-Deser-Misner decomposition of the metric as
\[
ds^2 = -N^2c^2dt^2 + g_{ij}(dx^i + N^i dt)(dx^j + N^j dt),
\]
where \( N \) is a lapse function, \( N^i \) is a shift vector, and \( g_{ij} \) is a three-dimensional spatial-metric. The HL gravity is formulated by an anisotropic scaling between time and space, \( t \rightarrow b^2 t \) and \( x^i \rightarrow b x^i \) and the action of \( z = 3 \) with the softly broken detailed balance condition is given by [6][10]

\[
I_{\text{HL}} = \int dt dx^i \sqrt{g} N \left[ \frac{2}{\kappa^2} (K_{ij} K^{ij} - \lambda K^2) - \frac{\kappa^2}{2c^4} (C_{ij} - \frac{\mu \kappa^2}{2} R_{ij}) (C^{ij} - \frac{\mu \kappa^2}{2} R^{ij}) + \frac{\kappa^2 \mu^2 (4\lambda - 1)}{32(3\lambda - 1)} \left( R^2 + 4(\omega - \Lambda_W) 4\lambda - 1 + 12\Lambda_W^2 \right) \right],
\]
which is sometimes called deformed HL gravity in the literature. Note that \( K_{ij} \) is the extrinsic curvature where \( \frac{\partial g_{ij}}{\partial t} C_{ij} \equiv \dot{g}_{ij}/c^2 \) denotes \( \dot{g}_{ij} \), \( C_{ij} \) \( i \neq j \) denotes \( \partial g_{ij}/\partial t \), \( C_{ij} \equiv \frac{\partial g_{ij}}{\partial t} \) is the Cotton-York tensor, and \( R_{ij} \) and \( R \) are the three-dimensional spatial Ricci tensor and Ricci scalar, respectively. Moreover, \( \kappa^2 \) is a coupling related to the Newton constant \( G_N \), and \( \lambda \) is an additional dimensionless coupling constant. Here \( \omega \) is essential to have asymptotically flat solutions, though it breaks the detailed balance condition. The coupling constants \( \mu \), \( \Lambda_W \), and \( \zeta \) come from the three-dimensional Euclidean topologically massive gravity action [29][30].

When \( \lambda = 1 \), the Einstein-Hilbert action can be recovered in the IR limit by identifying the fundamental constants with \( c = (\kappa^2/4)[\mu^2(\omega - \Lambda_W)/(3\lambda - 1)]^{1/2} \), \( G_N = \kappa^2 c^2/32\pi \), and \( \Lambda = -(3/2)\Lambda_W^2/(\omega - \Lambda_W) \), which represent the speed of light, the gravitational constant, and the cosmological constant, respectively.

The total action under consideration is given by \( I_{\text{tot}} = I_{\text{HL}} + I_{\text{mat}} \), where \( I_{\text{mat}} \) represents the matter action that will be specified without an explicit form by assuming the perfect fluid and choosing a EoS. In addition, we consider hereafter the asymptotically flat geometry, i.e. \( \Lambda = 0 \), for simplicity. Now if we consider a static, spherically symmetric metric ansatz,

\[
ds^2 = -e^{2\phi(r)} c^2 dt^2 + \frac{dr^2}{1 - f(r)} + r^2 (d\theta^2 + \sin^2 \theta d\phi^2),
\]
the equations of motion in HL gravity with the perfect fluid are given by

\[
\rho = \frac{c^2}{16\pi G_N r^2} \left( 2r\omega f + \frac{f^2}{r} \right),
\]

\[
p = \frac{c^4}{16\pi G_N r^4} \left[ f(f - 2r^2\omega) + 4r(1 - f)(f + r^2\omega) \Phi' \right],
\]

\[
p' = -(\rho c^2 + p) \Phi',
\]
where \( \rho \) and \( p \) are the mass density and the pressure of perfect fluid, respectively, and the prime denotes d/dr. Note that among the six parameters in the action (1), the only remaining free parameter is \( \omega \). To solve this equation, let us now replace \( f = -\sqrt{\omega} \sqrt{r^2\omega + 4G_N c^2 m} \), which, for \( m(r) = M \), approaches to \( f \rightarrow \infty \) \( 2G_N c^2 M/r - 2G_N c^4 m^{-1} M^2/r^4 + \ldots \), where the first term is nothing but the Schwarzschild solution and the following terms are HL-corrections depending upon \( \omega \) parameter. For the KS solution, we see \( e^{2\phi(r)} = 1 - f(r) \) while Eq. (30) governs the behavior of (tt)-component of the metric which is coupled to matter in the generic case. Note that the horizons of KS black hole are located at \( r_{\pm} = G_N c^{-2} \left( M \pm \sqrt{M^2 - M^2} \right) \), where \( M = (2G_N c^4 m^{-1})^{-1/2} \) is the critical mass of forming horizons.

With the Planck mass \( m_P \) and Planck length \( l_P \), \( M_c \) can be rewritten as \( M_c = m_P / \sqrt{2\omega l_P^2} \), which reduces to \( M_c = m_P / \sqrt{2} \) for \( \omega = \ell_P^2 \) and to \( M_c \approx M_\odot \) for \( \omega \approx 2.293 \times 10^{-11} \text{cm}^{-2} \). For a constant \( \omega \), the black hole horizon exists with \( M \geq M_c \), while a naked singularity appears when \( M < M_c \). The detailed behavior of forming naked singularities and wormholes in Hořava gravity is studied in Ref. [31]. On the other hand, for a fixed \( M \), one finds \( \omega \approx \omega_c \approx (2G_N c^4 M^2)^{-1/2} \) to form a black hole. For a solar mass KS black hole, its horizon exists only if \( \omega \geq \omega_c \approx 2.293 \times 10^{-11} \text{cm}^{-2} \). In quantum regime, however, KS geometry turns out to be regular [32], which may safely avoid the naked singularity problem.

Now, we can rewrite Eq. (3) in terms of \( m(r) \), \( p(r) \),
and $\rho(r)$ as

$$m' = 4\pi r^2 \rho,$$

$$p' = \frac{(\rho c^2 + p) r \omega [(1 + \tilde{\rho}) - \sqrt{1 + 4\tilde{\rho} - \tilde{p}}]}{\sqrt{1 + 4\tilde{\rho} + r^2 \omega (1 - \sqrt{1 + 4\tilde{\rho}})}},$$

(4a)

(4b)

where $\tilde{\rho} = G_N c^{-2} \omega^{-1} m r^{-3}$ and $\tilde{p} = 4\pi G_N c^{-4} \omega^{-1} p$.

Note that we can expand $p'$ in terms of $1/\omega \to 0$ using the relation $\sqrt{1 + \epsilon} \approx 1 + \epsilon/2 + \cdots$ for $\epsilon \ll 1$ as

$$p' \approx -\frac{G_N c^{-2} m r^{-2} (\rho c^2 + p) [1 + 4\pi c^{-2} r^3 p/m]}{1 - 2G_N c^{-2} m/r}$$

(5)

in the leading order, which reproduces the TOV equation in GR.

**Neutron Star Structure in HL gravity.** — For $\rho$ and $p$ in Eq. (4), various EoS models have been suggested. Amongst various models, we, in particular, select four EoS models which result the maximum mass of a neutron star to be $\geq 2M_\odot$ from conventional TOV calculation in GR and the results fit to the recent observations of $\sim 2M_\odot$ NS [25, 26]. They are APR4 [21], MPA1 [22], MS1 [23], and WFF1 [24]. MPA1 and MS1 are respectively derived by relativistic Brueckner-Hartree-Fock and relativistic mean field theory. APR4 and WFF1 are derived by variational method but with different nucleon potential models (see Ref. [21] for the difference in used models). All of these models assume that the nuclear matter is containing neutrons, protons, electrons, and muons. Additionally, in order to consider crust structure near the surface of a NS, we impose Skyrme-Lyon model [33] for the low-density region in addition to above models as shown in Ref. [34].

We adopt the 5th order Runge-Kutta method with 4th order error control [27, 28] to solve $m$ and $p$ in Eq. (4) simultaneously. And we impose the boundary conditions $\rho(0) = \rho_c$, $p(0) = p_c$ at the center $r = 0$ where $(\rho_c, p_c)$ satisfies a given EoS. By solving Eq. (4) numerically, we obtain a mass and pressure profile of the star which are $m(r)$ and $p(r)$, respectively. Then, the radius $R$ is determined by the starting point of zero pressure such that $p(r) = 0$ for $r \geq R$. And we get the mass by $M = m(R)$. Thereby, varying $\rho_c$, we produce the series of $(M, R)$, the mass-radius relation, which describes the static profile of spherically symmetric NS structure for a given EoS model. The exterior of the NS is naturally described by the KS vacuum solution $e^{2\Phi(r > R)} = 1 - f(r > R) = 1 + r^2 \omega - \sqrt{\omega (r^2 \omega + 4G_N c^{-2} M)}$ and an integration constant of the metric function $\Phi(r)$ is fixed by this boundary condition at $r = R$.

In order to compare the result in HL gravity to GR, the mass-radius relation in GR is depicted with star-marked maximum mass in Fig. 1. In HL gravity, the mass-radius relations are plotted by varying parameter $\omega$ from $2 \times 10^{-12}$ cm$^{-2}$ to $10^{-9}$ cm$^{-2}$. One finds that this range is far above of $\omega_{\text{min}} \sim 10^{-48} - 10^{-16}$ cm$^{-2}$ constrained from Refs. [15-17]. In this Letter, we empirically set this range since it is sufficient to see the modification of the HL theory to the structure of a NS. We also mark the mass and radius $(M_{\text{max}}, R_{\text{max}})$ of the heaviest NS in HL gravity for each EoS in Fig. 1. The behavior in HL gravity shows that $M_{\text{max}}$ increases as $\omega$ becomes small in all cases of selected EoS models and the radius of NS with a given mass, e.g. $1.4M_\odot$, as well. This implies that HL gravity with smaller $\omega$ has a larger deformation from GR. In addition, we place horizons of Schwarzschild and KS black holes at top-left corner of Fig. 1 to see that in both HL and GR, the radius of a stable NS is larger than the horizon radius of a black hole with the same mass as desired.
We now focus on the $\omega$-dependency in mass and radius of the heaviest NS in HL gravity as shown in Fig. 2. Each plot shows $\omega$-dependent profiles of $M_{\max}$ and $R_{\max}$ in HL and those in GR are represented as dashed lines for comparison. To find the relation, we fit a curve to the data obtained from each model. The relations are given by $M_{\max}^{\text{HL}} \approx M_{\max}^{\text{GR}} [1 + (A_1 \omega^{-1} + A_2 \omega^{-2})]$ and $R_{\max}^{\text{HL}} \approx R_{\max}^{\text{GR}} [1 + (B_1 \omega^{-1} + B_2 \omega^{-2})]$. The detailed fitting coefficients are summarized in the Table I.

Here we can see that HL gravity can reproduce GR when $\omega \gtrapprox 1.68 \times 10^{-9}$ or $1.68 \times 10^{-10}$ cm$^{-2}$ up to tolerances $|\Delta M_{\max}| (M_{\max}^{\text{GR}})^{-1}, |\Delta R_{\max}| (R_{\max}^{\text{GR}})^{-1} \lesssim 10^{-3}$ or $10^{-2}$, respectively, where $\Delta$ represents the difference between a quantity in HL from that in GR. So we see that TOV equation profile with $\omega \gtrapprox 1.68 \times 10^{-9}$ cm$^{-2}$ in HL gravity has the same behavior with that in GR for all considered EoS models. We also see that the smaller value of $\omega$ in HL gravity leads the larger deformation from GR.

Next, we look the $r$-dependent profiles at the interior of NS with APR4 case which gives the closest maximum mass in GR amongst tested models to the upper limit of the NS mass $^{[35]}$. The other EoSs show similar behavior. In the upper panel of Fig. 3a considering APR4, we plot mass profiles of the following NSs in HL with $\omega = 2 \times 10^{-12}$ cm$^{-2}$ and compare them with (i) the heaviest NS in GR; (ii) the heaviest NS in HL, (iii) a NS of the same mass $M_{\max}^{\text{GR}}$ as the NS in (i), and (iv) a NS of the same central density $\rho_c$ as the NS in (i). In the lower panel, the gravitational acceleration inside the NSs (i)-(iv) are plotted. The surface gravity of a NS is given by $g_{\text{NS}} = -(1 - f(R))^{-1/2} f'(R)/2$ $^{[36, 37]}$. Similarly, the gravitational acceleration that felt by an observer on a fixed $r$ inside a NS can be calculated as $a(r) = (1 - f(r))^{1/2} \Phi'(r)$, which naturally becomes $g_{\text{NS}}$ at the surface $R$. It is easily seen that the HL gravity is relatively weaker than the GR near the center of a NS. Even though the NS (ii) is much heavier than the NS (i), its gravity is weaker at $r \lesssim 6$ km, which explains why the NS heavier than the NS (i) remains stable in HL without collapsing into a black hole. The fact that HL gravity is
weaker than GR is easily seen by the curve of the same-mass NS (iii) in the bottom panel of Fig. 3a. Note that the maximum accelerations inside the heaviest NSs in GR and HL are almost the same; however, they may be slightly different according to the EoS models.

Meanwhile, the profiles of density and pressure inside the NS are also shown in the upper panel and lower panel, respectively, in Fig. 3b. Due to the weaker gravitational attraction inside the same-mass NS (iii), the pressure and density increase slowly as $r$ goes into zero and the central density becomes much smaller than the NS (i), which is consistent with larger radius of the NS (iii). The central density of NS (iv) is the same with that of the NS (i), and the density and pressure decrease slowly as $r$ increases and vanish at farther radius, so that the NS (iv) becomes naturally heavier than the NS (i). Note that the weaker gravitational attraction in HL gravity can be easily understood by computing Kretschmann curvature scalar defined by $\mathcal{K} = R_{\alpha\beta\gamma\delta}R^{\alpha\beta\gamma\delta}$. Indeed, the Kretschmann curvature scalar of KS solution can be compared to that of Schwarzschild solution with the same mass and radius. In GR, $\mathcal{K}_{\text{SBH}} = 48G_N^2c^{-4}M^2r^{-6}$ while $\mathcal{K}_{\text{KS}} \xrightarrow{r \to r_c} \frac{(81)}{4}G_NN^{-2}\omega^{-3}1 + \frac{O(r/r_c)^{3/2}}{2}$ and $\mathcal{K}_{\text{KS}} \xrightarrow{r \to r_c} \frac{1}{2} + \frac{O(r/r_c)^3}{3}$ in HL gravity, where $r_c \equiv (4G_Nc^{-2}\omega^{-1})^{1/3}$ is a characteristic scale of KS solution. The overall behavior of Kretschmann curvature scalar of KS solution in HL is slowly growing rather than Schwarzschild solution in GR, which implies that gravitational force induced by the curvature scalar in HL is weaker than that in GR in the limit of $r \to 0$. Then the mass obtained by integrating Eq. (4a) from $r = 0$ to $r = R$ also becomes heavier, eventually. However, the results of $M_{\text{max}} \gtrsim 3 M_\odot$ in HL gravity are placed not only beyond the upper limit of theoretical estimation on the upper limit [35] but far beyond the upper limit of the mass of NS [25, 26, 39–42].

Discussions.

Now it is plausible to note that the EoS models for the NS structure do not show significant differences in the structure of a white dwarf. It is mainly because they are almost the same except the physics at the nucleus size. Likewise, the HL gravity deforms the UV behavior and does not change the structure of a white dwarf significantly. The behavior of white dwarfs as well as NSs in GR and HL gravity are depicted in Fig. 4. The relativistic EoS for white dwarfs is obtained by assuming that the dominant chemical components of white dwarfs are $^4\text{He}$, $^{12}\text{C}$ or $^{16}\text{O}$, that is the ratio between the atomic mass number $A$ and the number of electrons $Z$ is given by $A/Z = 2$ [43]. The EoS models for NSs do not reflect the Chandrasekhar limit correctly, since they are obtained by considering the physical situations beyond the degeneracy pressure of electrons such as neutronization.

On the other hand, the conventional measurement on the mass and radius of a NS is done with observation on a pulsar, rotating NS, consisting a binary system. Depending on the type of pulsar, different approaches have been implemented to measure or to constraint the mass and radius (for more details, see Ref. [44] and references therein). Additionally, from a recent detection of GW170817, radiated from a merger of binary NS system, now it is possible to consider gravitational-wave (GW) as an additional approach to estimate the mass and radius of NS [33–45]. Hence, in order to discuss the viability of HL gravity with physical observables of NS, we need to address some aspects of HL gravity more rigorously such as (i) whether it can pass the tests on post-Keplerian parameters [44], (ii) whether it is possible to derive GW waveforms with HL gravity to perform the parameter estimation as done with GR, (iii) how the tidal deformability of NS in HL gravity differs from that in GR, and so on.

In conclusion, it is shown that the HL gravity reveals a deformed feature at short distance regime due to its relatively weaker gravitational force than that of GR. We see that the deformation from GR is sensitive to the choice of parameter $\omega$. In addition, the parameter $\omega$ is determined from the lower limit of the current observational results obtained from GR or post-Keplerian. For the validation of HL gravity, theoretical investigations and future observations of GWs from compact binary mergers containing at least one NS will give more constraints on the physical observables of NS and, eventually, will determine the fate of HL gravity.

JJO would like to thank Parada Hutauruk for helpful discussions at the early stage of this work. The authors would like thank Chang-Hwan Lee and Young-Min Kim for fruitful discussions on neutron stars and tidal deformations related to GW170817 and also thank Olivier Mi-
nazzoli, Nathan Johnson-McDaniel, Noah Sennett, Harald Pfeiffer, Tjonnie Li, and Chris Van Den Broeck for helpful comments and discussion on this work. KK was partially supported by a grant from the Research Grants Council of the Hong Kong (Project No. CUHK 14310816 and CUHK 24304317) and the Direct Grant for Research from the Research Committee of the Chinese University of Hong Kong. CP was partially supported by Basic Science Research Program through the National Research Foundation of Korea (NRF) funded by the Ministry of Education (NRF-2018R1D1A1B07041004).

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