Do massive compact objects without event horizon exist in infinite derivative gravity?

Alexey S. Koshelev$^{1,2}$ and Anupam Mazumdar$^3$

$^1$Departamento de Física and Centro de Matemática e Aplicações, Universidade da Beira Interior, 6200 Covilhã, Portugal
$^2$Theoretische Natuurkunde, Vrije Universiteit Brussel and The International Solvay Institutes, Pleinlaan 2, B-1050, Brussels, Belgium
$^3$Van Swinderen Institute, University of Groningen, 9747 AG, Groningen, The Netherlands

Einstein’s General theory of relativity is plagued by cosmological and blackhole type singularities. Recently, it has been shown that infinite derivative, ghost free, gravity can yield non-singular cosmological and mini-blackhole solutions. In particular, the theory possesses a mass-gap determined by the scale of new physics. This paper provides a plausible argument, not a no-go theorem, based on the Area-law of gravitational entropy that within infinite derivative, ghost free, gravity non singular compact objects in the static limit need not have horizons.

I. INTRODUCTION

Einstein’s theory of general relativity (GR) is extremely successful in the infrared (IR) $^1$, matching all the current observations including predictions of gravitational waves from coalescing binary blackholes $^2$. The only drawback arises in its predictions in the ultraviolet (UV) regime, where there exists classical singularities, such as blackhole and cosmological Big Bang singularities, see $^3$. There exists a vicious circle in GR, which inevitably leads to a collapse of a normal matter satisfying all the known energy conditions $^4$.

Since in GR both energy and pressure gravitates, therefore, any normal matter satisfying strong, weak and null energy conditions will always lead to a focusing of either time-like and/or null rays according to the Raychouhdhury’s equation $^5$, which will lead to a formation of a trapped surface, and an apparent horizon, see $^6$. In the static limit, both apparent and event horizon coincides, and the metric potential is given by the Schwarzschild metric, which asymptotes to the Minkowski far away from the source. The potential is denoted by $\Phi \sim \frac{Gm}{r}$, where $m$ is the mass of a blackhole, and $G$ is the Newton’s constant $^7$. In the Schwarzschild’s geometry, where $r \leq 2Gm$, the metric potential keeps growing all the way to $\Phi \rightarrow \infty$ as $r \rightarrow 0$.

In the cosmological context, in the homogeneous and isotropic Universe, although the origin of singularity is slightly different than that of a blackhole, but there is one common ingredient, that is the energy density blows up as $t \rightarrow 0$, and so is the metric potential.

The aim of this paper is to break this vicious circle of inevitable collapse of matter in gravity without violating any of the energy conditions. Our arguments will be based on a static case in this paper, so strictly speaking they will not be directly applicable to a cosmological setup. There are two ways we can try to avoid inevitable singularity in a gravitational theory: (a) First option is to make gravity repulsive at short distances and small time scales, in the UV, such that a delicate balance of matter and gravitational pressure would halt the collapse problem, and therefore avoiding singularity. This approach will inevitably lead to introduction of ghosts in the gravitational sector. (b) The second option is to make gravitational interaction sufficiently weak, such that in the UV the spacetime becomes regular and the gravitational force between particles vanishes, $F_g \rightarrow 0$, therefore the gravitational binding energy ceases to be singular within a finite length and time scale. We will follow the second strategy here.

Here we will not attempt to provide a mathematical proof for avoiding either singularity or event horizon for systems as massive as astrophysical solar massive objects, but will present some arguments of plausibility regarding the possible static properties of non-singular compact objects.

II. GHOST FREE INFINITE DERIVATIVE GRAVITY

We wish to seek a theory of gravity, which allows weak field limit in the entire region of spacetime, i.e. $2\Phi < 1$ for all values of $r$. In the static and spherically symmetric geometry, this would mean $\Phi < 1$ for both $r \geq 2Gm$ and $r \leq 2Gm$. Indeed, at sufficiently large values, we would like to recover $1/r$-fall of Newtonian potential. Therefore, the desirable space-time metric should be the linearized form always, specifically in a static geometry:

$$ds^2 = -(1 - 2\Phi)dt^2 + (1 + 2\Phi)dr^2,$$

where $dr^2 = dx^2 + dy^2 + dz^2$, and $2\Phi < 1$, such that the solution remains perturbative for a given modification of GR.

It has been shown recently that infinite derivative gravity (IDG) can avoid cosmological and blackhole singularities by making the gravitational action ghost free $^8$, such that the propagating degrees of freedom remain massless spin-2 and spin-0 components $^9$. In addition to this, IDG also allows massless spin-2 and one

---

$^1$ We will also use the 4 dimensional gravitational Planck constant: $M_p = 1/\sqrt{\mathcal{G}} \sim 10^{19}$ GeV.
massive spin-0 components to propagate, analogous to the Brans-Dicke gravity, see \[12,13\]. The most general ghost free IDG action can be recast as \[8,13\]:

\[
S = \int d^4x \sqrt{-g} \left[ M_p^2 R + R \mathcal{F}_1(\Box) R + R^{\mu\nu} \mathcal{F}_2(\Box) R_{\mu\nu} + R^{\mu\nu\lambda\sigma} \mathcal{F}_3(\Box) R_{\mu\nu\lambda\sigma} \right] (2)
\]

where \( \Box = \Box / M_p^2 \), and \( M_s \) is the new scale of physics which appears below the 4 dimensional Planck mass, i.e. \( M_s \leq M_p \). It has been shown in Refs. \[10,14–18\] that \( M_s \) harbours the scale of non-locality at the level of quantum interactions.

The three gravitational form-factors behave as \( 2 \mathcal{F}_1 + \mathcal{F}_2 + 2 \mathcal{F}_3 = 0 \) around a Minkowski background \[8\], in order to propagate only the massless graviton. In order to avoid any new dynamical degrees of freedom, that includes tachyons and/or ghosts, the propagator must be suppressed by exponential of an entire function \[8,10\].

\[
\Pi(k^2) \sim e^{-\gamma(k^2)} \left[ P^{(2)} - P^{(0)} \right] \frac{k^2}{2k^2}, \quad (3)
\]

where \( P^{(2)} \), \( P^{(0)} \) are spin-projection operators respectively. The exponential factor with \( \gamma(k^2) \) being an entire function is mathematically an unique option in order to avoid new poles at finite momenta. Surely, \( \gamma(k^2) \) must obey certain conditions. In particular, it should grow at large momenta such that \( \gamma(k^2) \) decays in the UV, signalling the weakening of the graviton propagation for physical degrees of freedom in UV, and in the limit when \( k \to 0 \), in the IR the propagator matches exactly the behaviour of Einstein’s gravity. The simplest examples of such an entire functions are polynomials, and we concentrate our discussions with \( \gamma(k^2) = k^2 / M_s^2 \). Note that in \[23\] other polynomials and generic series for function \( \gamma(k^2) \) were analyzed and proven to be compatible with the theory. However, the main idea of the current paper can be presented by using just the simplest monomial, and moreover, conceptually our conclusions will not be affected by considering a more general form of \( \gamma(k^2) \).

In the limit when \( M_s \to \infty \), the theory comes back to the predictions of the IR, i.e. the Einstein-Hilbert action \[4\]. The above modified propagator ensures that in the UV, the propagator is exponentially suppressed, thus improving upon the UV behaviour of gravity. The vertex operator on the other hand gets exponential enhancement, therefore at a quantum level, the interactions become non-local. The superficial degree of divergence suggests that such gravitational theory becomes power counting renormalizable for loop \( L > 1 \) \[11, 18\].

III. GRAVITATIONAL POTENTIAL FROM A SOURCE

A very interesting consequence of this suppressed propagator can be illustrated by the fact that for a point-source, the linearized equations of motion yields \[8\]:

\[
e^{-\Box / M_s^2} \Box \Phi = 8\pi G \rho = m \delta^3(r). \quad (4)
\]

In the static limit, when \( \Box = \nabla^2 \) the gravitational potential can avoid \( 1/r \) singularity, which is impossible to avoid otherwise in GR. The metric potential has a solution given by \[8\]:

\[
\Phi(r) = \sqrt{\frac{\pi}{2 M_s^3}} \text{erf}(M_s r/2), \quad (5)
\]

where for \( r \gg M_s \), the potential falls as the Newtonian limit: \( 1/r \), and for \( r < 2/M_s \), the potential asymptotes to a constant value, and the gravitational force vanishes:

\[
\Phi \sim m M_s / M_p^2 < 1, \quad F_g \to 0, \quad \text{for } r < \frac{2}{M_s}, \quad (6)
\]

and for mass enclosed within such a radius follows:

\[
m < M_p^2 / M_s. \quad \text{For } r > 2/M_s, \quad \text{the erf} \to 1, \quad \text{and therefore one recovers the standard } 1/r\text{-fall of gravity. As it has been shown earlier in } 25,26, \quad \text{a constant asymptotic in the region } r < 2/M_s \text{ and standard } 1/r\text{ fall-off at large distances are the universal features of IDG, while the value of the cross-over scale } 2/M_s \text{ depends on the particular form of function } \gamma(k^2) \text{ in Eq. (6).}
\]

Namely, the factor of 2 in the cross-over scale is subject to the lowest degree of \( k^2 \) in the series expansion of \( \gamma(k^2) \). For, \( \gamma(k^2) \approx (k^2 / M_s^2)^n + O(k^{2n+2}) \) with \( n > 1 \) the factor of 2 will be subject to modification. In spite of these details, let us stick with the lowest order polynomial for the purpose of our discussion.

In this regard the IDG possess a scale, known as mass gap, determined by the scale of non-locality \( M_s \) \[8,23\].

The current constraint on \( M_s \) arises from table-top laboratory experiment, which has seen no deviation from Newtonian gravity up to \( 5.6 \times 10^{-5} \) m \[20\]. This limit translates to: \( M_s \geq 0.004 \text{ eV} \) \[23\], and in order for \( \Phi < 1 \), we obtain a bound on mass

\[
m \leq \frac{M_p^2}{M_s} \sim 10^{25} \text{ grams}, \quad (7)
\]

\[3\] This has now been verified by Valeri Frolov and his collaborators in a dynamical context as well, see Refs. \[26,28\]. Note however that there are results \[44\] which contradict the results even at the linearized limit \[8,22\]. These papers \[44\] do not attempt to solve the complete equations of motion for the ghost free IDG, the complete equations of motion have been derived in Ref. \[24\]. There is another difference, in our case we have an explicit source term, Eq. (4), which is apparently lacking in their solution. It has been assumed that a vacuum solution will be similar to the Schwarzschild, but there is no explicit proof given.

---

Such exponential suppression in the propagator is quite common in theories with infinite derivatives, whether they arise in a bosonic sector, or in a fermionic sector \[18,20\]. In fact, such a propagator also arises in string field theory \[21,22\].
which is roughly the mass of the Moon. Therefore, for a Moon-like massive system would never generate a metric singularity at \( r = 0 \), and its potential will be regular all the way without forming any event horizon, i.e. \( 2\Phi < 1 \). If we had chosen \( M_s \sim M_p \), the bound on \( m \) would be:

\[
m \leq M_p \sim 10^{-5} \text{ grams},
\]

ideal for a planckian size object, i.e. \( r \sim 2M_p^{-1} \), where the metric potential is constant. Such regions of space-time, with a constant metric potential, could be treated as “plaquette”, which we denote here by \( \mathcal{U} \).

The property of a plaquette is solely determined by \( M_s \) and \( M_p \) within IDG. If we had chosen \( M_s = 10^{16} \text{ GeV} \), then such a plaquette would be able to hold \( m \leq M_p^2/M_s \sim 10^{-2} \) grams, without forming a singularity within \( r \leq 2/M_s \). Strictly speaking, this notion holds true for a static geometry and some of these arguments will modify in a time dependent case

### IV. NON-SINGULAR COMPACT OBJECT: SPACE-TIME FILLING PLAQUETTES

The key point to note here is that the plaquette’s mass is determined by the mass gap \( M_s \). Since the planckian energy density is the largest one can achieve, there exists a simple bound on a plaquette with a constant potential from Eq. (8):

\[
m^2 M_s^2 \leq M_p^4 \quad \text{for} \quad r < \frac{2}{M_s}.
\]

In the static limit, even if there are \( N \) such plaquettes, within a compact region of space-time, one would naturally expect the largest energy density to be still given by the planckian one:

\[
N^2 m^2 \left( \frac{M_s}{\sqrt{N}} \right)^2 \equiv M_{ns}^2 \frac{M_{eff}^2}{M_p^2} \leq M_p^4,
\]

for \( r < r_s = \frac{2}{M_{eff}} \),

where the rest mass of a non-singular compact object (NSCO) is now given by: \( M_{ns} = Nm \). Note that in order for the above condition to hold true, the effective non-local scale has shifted from \( M_s \) to a much lower value:

\[
M_{eff} \sim \frac{M_s}{\sqrt{N}}.
\]

A simple but compelling argument can be presented from the entropy point of view. Any gravitationally bound system is known to possess a gravitational entropy [34–37].

Let us imagine that the gravitationally bound space-time, \( \mathcal{M} \), is filled up with such constant potential plaquettes, \( \mathcal{U} \), such that the gravitational potential is constant over \( \mathcal{M} \) on a macroscopic scale, i.e. \( r_s = 2/M_{eff} \). If every plaquette, \( \mathcal{U} \), of size \( \sim 2/M_s \) denotes 1-unit of entropy, then the space-time filling \( N \) such constant potential plaquettes will have entropy scaled by \( N \)-units. One would expect similar result from computing the gravitational entropy for a gravitationally bound system, \( \mathcal{M} \), with a constant potential \( \Phi < 1 \) inside \( r \leq r_s \).

Note that the leading order contribution to the gravitational entropy is given by the Area-law, where \( S_g = \text{Area}/4G \), where Area is the enclosed area of such a bound system, and \( G \) is the Newton’s constant [34–37]. At the leading order the Area-law holds true for the ghost free IDG as well, see [38]. As long as the metric potential, \( \Phi < 1 \), the Wald’s gravitational entropy comes out to be proportional to the Area for the static and spherically symmetric bound system.

Therefore, in our case, the gravitational entropy for \( \mathcal{M} \) would scale as \( \propto 4\pi r_s^2 \), where \( r_s = \sqrt{N}/M_s \):

\[
S_g = \frac{\text{Area}}{4G} \propto N \left( \frac{M_p}{M_s} \right)^2 \sim \left( \frac{M_p}{M_{eff}} \right)^2.
\]

This analysis suggests that for a system with a mass-gap, there is a possibility to shift the scale of non-locality from \( M_s \) to \( M_{eff} = \sqrt{N}/M_s \) for a gravitationally bound system, i.e. \( \Phi \sim \text{constant within radius} \ r_s = 2/\sqrt{M_{eff}} \).

Based on this analysis we can now ask; could we form a super-massive NSCO with a radius bigger than the Schwarzschild’s radius, \( r_{sch} \), i.e. \( r_s \geq r_{sch} \). Within IDG it is indeed possible to construct \( \mathcal{M} \), which has a constant metric potential. In fact, following the arguments of the last section, we can now find the metric potential to be:

\[
\Phi(r) = \sqrt{\frac{\pi}{2}} \frac{M_{ns}}{M_p^2} \text{erf}(r M_{eff}/2),
\]

where \( M_{ns} = Nm \) is the mass of NSCO. The solutions of the above equation yields:

\[
\Phi \sim \begin{cases} \frac{M_s M_{eff}}{M_p^2} < 1, & r < r_s = \frac{2}{M_{eff}}, \\ \frac{GM_{ns}}{r} < 1, & r > r_s = \frac{2}{M_{eff}}. \end{cases}
\]

Now, we can seek under what conditions \( r_s \geq r_{sch} \):

\[
r_s = \frac{2}{M_{eff}} \geq r_{sch} = \frac{2 M_{ns}}{M_p^2} \quad \Rightarrow \quad M_{ns}^2 M_{eff}^2 \leq M_p^4,
\]

5 This scaling of non-locality can also be seen as a property of field theory with infinite derivatives at a quantum level. It has been shown that the scalar counterpart of Eq. (4), which preserves the scaling and the shift symmetry, see [38], has a similar property in the scattering amplitudes. The scattering amplitude is exponentially suppressed in the UV, when centre of mass energy exceeds \( M_s \). However, when multiple scatterings are considered, the scattering amplitude becomes suppressed by the new effective scale \( M_{eff} \rightarrow M_s/\sqrt{N} \) for large \( N \)-limit [39].
which is a similar condition as that of the scaling of the entropy argument, see Eq. (12), and also Eq. (10).

We can also estimate what should be the value of $N$ if NSCO were made up of a billion times the solar mass, i.e. $10^{12} \times 10^{33} \sim 10^{45}$ grams. For $M_*=M_p$, $N \geq 10^{100}$, while, for $M_* = 10^{16}$ GeV, $N \geq 10^{94}$, and for $M_* = 10^{4}$ GeV, $N \geq 10^{70}$. Indeed, a NSCO can hold a large $N$, which signifies large amount of entropy, very similar to the case of a typical blackhole within GR.

V. ABSENCE OF AN EVENT HORIZON

One of the properties of NSCO within IDG is that the absence of an event horizon, since $2\Phi < 1$ in the macroscopic region of space-time, as large as that of the Schwarzschild’s radius, $r_s \geq r_{sch}$, see Eq. (13). By construction the metric potential is given by Eq. (14), in conjunction with Eq. (15).

Indeed Sagitarus A∗, which harbours a huge mass in the centre of our Milkyway [41], yields $1/r$ gravitational potential for a distant observer, $r > r_s > r_{sch}$. In the idealised case, assuming in the static scenario, when the observer comes close to the vicinity of its Schwarzschild radius the gravitational potential tends to become constant for $r < r_{sch} < r_s$. The transition is indeed smooth and determined by the full solution, Eq. (16).

The absence of an event horizon is indeed a very compelling reason why such NSCO should be further studied. The Hawking’s information-loss paradox [42] can be resolved amicably in the absence of any stretched horizon. The absence of an event horizon means that the information can never get lost, but indeed, NSCO can modify how the information can be retrieved or scrambled due to large $N$ plaquettes, which requires further study, see for instance [43]. Here, we wish to draw some similarity with the fuzz-ball scenario [45], where the physical situation has some common feature, though the intricate details are very different. In the fuzz-ball scenario there is no event horizon either, and it also reiterates that the scale of quantum gravity need not be localised on a small scale, but it can be enlarged to a macroscopic length scale in a self-gravitating system [46]. This happens purely due to stringy origin, i.e. presence of non-local objects such as winding modes, branes, etc. In our case, higher derivatives are akin to $\alpha'$ corrections in string theory, therefore there might be some interesting connection which one might be able to exploit between these two seemingly different approaches. Furthermore, besides fuzz-ball scenario, there are examples such as gravastars [46,47], and blackstars [48,51], which mimic blackholes without event horizons.

VI. CONCLUSION

In the Einstein-Hilbert gravity, it is a challenge to avoid $1/r$ singularity in the metric potential for a static and spherically symmetric background. The IDG provides a compelling candidate of quantum gravity where the gravitational potential, $\Phi$, can be made constant in the UV, such that at short distances the gravitational force vanishes. The scale is typically determined by the mass gap, i.e. the scale of non-locality $M_*$.

VII. ACKNOWLEDGEMENT

The authors would like to thank Tirthabir Biswas, Tomi Koivisto, Valeri Frolov, Gaetano Lambiase, Terry Tomboulis and Shinji Mukohyama for discussing various aspects of this problem. The authors would also like to thank the anonymous referee for a very helpful discussion, and for numerous suggestion to help improving the quality of the paper. AK is supported by FCT Portugal investigator project IF/01607/2015, FCT Portugal fellowship SFRH/BPD/105212/2014 and in part by FCT.
[1] C. M. Will, “The Confrontation between General Relativity and Experiment,” Living Rev. Rel. 17, 4 (2014) doi:10.12942/lrr-2014-4 [arXiv:1407.7377 [gr-qc]].

[2] B. P. Abbott et al. [LIGO Scientific and Virgo Collaborations], “Observation of Gravitational Waves from a Binary Black Hole Merger,” Phys. Rev. Lett. 116 (2016) no.6, 061102 doi:10.1103/PhysRevLett.116.061102 [arXiv:1602.03837 [gr-qc]].

[3] S. W. Hawking and G. F. R. Ellis, “The Large Scale Structure of Space-Time,”

[4] R. M. Wald, “General Relativity,” Chicago, Usa: Univ. Pr. (1984) 491p

[5] R. Penrose, “Gravitational collapse and space-time singularities,” Phys. Rev. Lett. 14, 57 (1965). doi:10.1103/PhysRevLett.14.57 S. W. Hawking and R. Penrose, “The Singularities of gravitational collapse and cosmology,” Proc. Roy. Soc. Lond. A 314, 329 (1970). doi:10.1098/rspa.1970.0021 S. Hawking, “The occurrence of singularities in cosmology. III. Causality and singularities,” Proc. Roy. Soc. Lond. A 300, 187 (1967).

[6] A. Raychaudhuri, “Relativistic cosmology. 1.,” Phys. Rev. 98, 1123 (1955). doi:10.1103/PhysRev.98.1123

[7] V. Frolov and A. Zelnikov. 2011. Introduction to black hole physics. Oxford University Press.

[8] T. Biswas, E. Gerwick, T. Koivisto and A. Mazumdar, “Towards singularity and ghost free theories of gravity,” Phys. Rev. Lett. 108 (2012) 031101 doi:10.1103/PhysRevLett.108.031101.

[9] T. Biswas, A. Mazumdar and W. Siegel, “Bouncing universes in string-inspired gravity,” JCAP 0603 (2006) 009 [hep-th/0508194].

[10] E. Tomboulis, Phys. Lett. B 97, 77 (1980). E. T. Tomboulis, Renormalization And Asymptotic Freedom In Quantum Gravity, In "Christensen, S.m. ( Ed.): Quantum Theory Of Gravity*, 251-266. E. T. Tomboulis, Superrenormalizable gauge and gravitational theories, hep-th/9702146.

[11] L. Modesto, Phys. Rev. D 86, 044005 (2012)

[12] T. Biswas, T. Koivisto and A. Mazumdar, “Nonlocal theories of gravity: the flat space propagator,” arXiv:1302.0532 [gr-qc].

[13] T. Biswas, A. S. Koshelev and A. Mazumdar, “Consistent higher derivative gravitational theories with stable de Sitter and anti-de Sitter backgrounds,” Phys. Rev. D 95 (2017) no.4, 043533 doi:10.1103/PhysRevD.95.043533 [arXiv:1606.01250 [gr-qc]]. T. Biswas, A. S. Koshelev and A. Mazumdar, “Gravitational theories with stable (anti-)de Sitter backgrounds,” Fundam. Theor. Phys. 183 (2016) 97 doi:10.1007/978-3-319-31299-65 [arXiv:1602.08175 [hep-th]].

[14] A. A. Tseytlin, “On singularities of spherically symmetric backgrounds in string theory,” Phys. Lett. B 363, 223 (1995) doi:10.1016/0370-2693(95)01228-7 [hep-th/9509050].

[15] W. Siegel, Stringy gravity at short distances, hep-th/0309093

[16] T. Biswas and N. Okada, “Towards LHC physics with nonlocal Standard Model,” Nucl. Phys. B 898 (2015) 113 doi:10.1016/j.nuclphysb.2015.06.023 [arXiv:1407.3331 [hep-ph]].

[17] T. Biswas, M. Grisaru and W. Siegel, “Linear Regge trajectories from worldsheet lattice partition field theory,” Nucl. Phys. B 708 (2005) 317 doi:10.1016/j.nuclphysb.2004.11.004 [hep-th/0409089].

[18] S. Talaganis, T. Biswas and A. Mazumdar, “Towards understanding the ultraviolet behavior of quantum loops in infinite-derivative theories of gravity,” Class. Quant. Grav. 32, no. 21, 215017 (2015) doi:10.1088/0264-9381/32/21/215017 [arXiv:1412.3367 [hep-th]].

[19] G. V. Efimov, “Non-local quantum theory of the scalar field,” Commun. Math. Phys. 5 (1967) no.1, 42. doi:10.1007/BF01664637 G. V. Efimov and S. Z. Seltser, “Gauge invariant nonlocal theory of the weak interactions,” Annals Phys. 67 (1971) no.1, 124. doi:10.1016/0003-4916(71)90007-8 V. A. Alebastrov and G. V. Efimov, “A proof of the unitarity of S-matrix in a nonlocal quantum field theory,” Commun. Math. Phys. 31 (1973) no.1, 1. doi:10.1007/BF01645588

[20] D. Evens, J. W. Moffat, G. Kleppe and R. P. Woodard, “Nonlocal regularizations of gauge theories,” Phys. Rev. D 43 (1991) no.2, 499. doi:10.1103/PhysRevD.43.499

[21] E. Witten, “Noncommutative Geometry and String Field Theory,” Nucl. Phys. B 268 (1986) 253. doi:10.1016/0550-3213(86)90155-0

[22] P. G. Freund and M. Olson, “Nonarchimedean Strings,” Phys. Lett. B 199 (1987) 186. doi:10.1016/0370-2693(87)91356-6

[23] P. H. Frampton and Y. Okada, “Effective Scalar Field Theory of F−adic String,” Phys. Rev. D 37 (1988) 3077. doi:10.1103/PhysRevD.37.3077

[24] W. Siegel, “Introduction to string field theory,” Adv. Ser. Math. Phys. 8 (1988) 1 [hep-th/0107094].

[25] J. Edholm, A. S. Koshelev and A. Mazumdar, “Behavior of the Newtonian potential for ghost-free gravity and singularity-free gravity,” Phys. Rev. D 94 (2016) no.10, 104033 doi:10.1103/PhysRevD.94.104033 [arXiv:1604.01989 [gr-qc]].

[26] V. P. Frolov, A. Zelnikov and T. de Paula Netto, “Spherical collapse of small masses in the ghost-free gravity,” JHEP 1506, 107 (2015) doi:10.1007/JHEP06(2015)107 [arXiv:1504.00412 [hep-th]]. V. P. Frolov and A. Zelnikov, “Head-on collision of ultra-relativistic particles in ghost-free theories of gravity,” Phys. Rev. D 93 (2016) no.6, 064048 [arXiv:1509.03356 [hep-th]].

[27] V. P. Frolov, “Mass-gap for black hole formation in higher derivative and ghost free gravity,” Phys. Rev. Lett. 115, no. 5, 051102 (2015) doi:10.1103/PhysRevLett.115.051102 [arXiv:1505.00492 [hep-th]].

[28] V. P. Frolov, “Notes on non-singular models of black holes,” Phys. Rev. D 94 (2016) no.10, 104056 doi:10.1103/PhysRevD.94.104056

[29] T. Biswas, A. Conroy, A. S. Koshelev and A. Mazumdar, “Generalized ghost-free quadratic curvature gravity,” Class. Quant. Grav. 31 (2014) 015022 Erratum: [Class. Quant. Grav. 31 (2014) 159501] doi:10.1088/0264-
