

Research Article

Probable Entropic Nature of Gravity in Ultraviolet and Infrared Limits—Part I: An Ultraviolet Case

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This work presents a study of the possibility for extending the well-known results of E. Verlinde concerning the entropic nature of gravity to the ultraviolet region (Planck's energies) and also the derivation of quantum corrections to Einstein Equations.

1. Introduction

In the last 15–20 years, new and very interesting approaches to gravity studies have been proposed, which may be divided into “thermodynamical” and “theoretical informational” approaches. The approach suggested in the pioneer work by Jacobson [1] has been considerably extended in a series of remarkable papers by Padmanabhan [2–13]. The paper by Verlinde [14] stating a secondary character of the gravitational interaction and its entropic nature was published in 2010 after the appearance in the ArXiv. The paper [14] introducing such specific terms as “entropic force” has been followed by numerous studies (e.g., [15–26] and others).

In this work, the author studies the possibility for extension of the results given in [14] to the ultraviolet region (Planck's energies) and presents the derivation of quantum corrections to Einstein Equations using the dimensionless small parameter α introduced by the author in his previous works [27–39].

2. Fundamental Quantities and Their High-Energy Deformation

In this section, the author uses the “ideology,” terms, and notation introduced in [14] to extend the corresponding results to the ultraviolet and infrared gravity regions. It may be stated that the results in [14] have been obtained for

the “medium” energies, that is, for the range of well-known energies, where the General Relativity (GR) is valid. But owing to modern knowledge, in the ultraviolet and infrared limits, gravity may be modified. As regards to the ultraviolet (Planck) scale, this idea has been proposed long ago [40–45], and in this situation we have to replace the word “may” by the word “must”. As for low energies, there are many recent publications considering the infrared (for great distances) modification of gravity (e.g., [46, 47]). And the modification can have a solid experimental status in the nearest future [48].

Naturally, when we are concerned with extension of some results to higher or lower energies, the principle of conformity must be executed undeviatingly; ongoing to the known energy scales, the known results must be reproduced.

In this section, it is shown that the fundamental quantities A , N , and T , defined in [14] and associated with the holographic screen (where A —its surface area, N —number of data bits “existing” on A , T —its temperature), may be supplementary defined for the region of high energies so that at normal energies they be coincident with values given in [14]. In the process, we take the corresponding quantities for the stationary Schwarzschild black hole as the most natural holographic object.

The idea is as follows: formulae for the correction of the fundamental quantities within the Generalized Uncertainty Principle (GUP) for the holographic screen \mathcal{S} from [14] are similar to those for the black hole.

Let us consider Section 5.2 (Derivation of Einstein Equations) in [14]. In this section, formula (5.32) for a ‘‘bit density’’ on the holographic screen is given as

$$dN = \frac{dA}{G\hbar}. \quad (1)$$

However, when the holographic principle [49–53] is valid, N is actually the entropy S up to the factor $S \sim N$, and hence from (1) it follows directly that

$$dS \sim \frac{dA}{G\hbar}. \quad (2)$$

What are the changes in S ongoing to high (Planck) energies? The answer to this question is already known owing to the fact that at these energies the Heisenberg Uncertainty Principle (HUP) is replaced by GUP [54–61]:

$$\Delta x \geq \frac{\hbar}{\Delta p} + \ell^2 \frac{\Delta p}{\hbar}, \quad (3)$$

where $\ell^2 = \alpha' l_p^2$ and α' is the dimensionless numerical factor. The well-known Bekenstein-Hawking formula for the black hole entropy in the semiclassical approximation [62, 63],

$$S^{\text{BH}} = \frac{A}{4l_p^2}, \quad (4)$$

is modified by the corresponding quantum corrections ongoing from HUP to GUP [64–67].

In particular, [65],

$$S_{\text{GUP}}^{\text{BH}} = \frac{A}{4l_p^2} - \frac{\pi\alpha'^2}{4} \ln\left(\frac{A}{4l_p^2}\right) + \sum_{n=1}^{\infty} c_n \left(\frac{A}{4l_p^2}\right)^{-n} + \text{const.}, \quad (5)$$

where the expansion coefficients $c_n \propto \alpha'^{2(n+1)}$ can always be computed to any desired order of accuracy.

The general form of quantum corrections for the black hole entropy derived in (5) remains valid for any horizon spaces and, in particular, for the holographic screen \mathcal{S} from [14]. Specifically, in [68, 69] the logarithmic correction was obtained in the following form:

$$S_{\text{ln}} = \frac{A}{4l_p^2} + \frac{\tilde{\alpha}}{4} \ln\left(\frac{A}{l_p^2}\right), \quad (6)$$

omitting the Boltzmann constant k_B as a factor and assuming it to be equal to unity in what follows. Higher-order corrections may be derived using the Taylor-series expansion in terms of the small parameter l_p^2/A :

$$S_{\text{GUP}} = \frac{A}{4l_p^2} + \frac{\tilde{\alpha}}{4} \ln\left(\frac{A}{l_p^2}\right) + \sum_{n=1}^{\infty} \tilde{c}_n \left(\frac{A}{l_p^2}\right)^{-n} + \text{const.} \quad (7)$$

in a similar way to the Taylor-series expansion of the right-hand side in (9) in terms of the small parameter $4l_p^2/A$. This is

valid as GUP gives the ultraviolet cutoff at the level of $l_{\text{min}} \sim l_p$.

In this way, at high energies we have $S \rightarrow S_{\text{GUP}}$, and hence $N \rightarrow N_{\text{GUP}}$. Assuming in the notation of [14] that

$$S = \frac{1}{4}N, \quad (8)$$

we directly obtain

$$N_{\text{GUP}} = \frac{A}{l_p^2} + \tilde{\alpha} \ln\left(\frac{A}{l_p^2}\right) + 4 \sum_{n=1}^{\infty} \tilde{c}_n \left(\frac{A}{l_p^2}\right)^{-n} + \text{const.} \quad (9)$$

Now, coming back to [14], in terms of N_{GUP} we can define the holographic screen area, as measured at high energies, by

$$A_{\text{GUP}} \equiv G\hbar N_{\text{GUP}}, \quad (10)$$

where G and \hbar are gravitational and Planck constants, respectively, and N_{GUP} is given by (9). Considering that we, similar to [14], assume that the speed of light $c = 1$, then, according to $l_p^2 = G\hbar$, from (9) and (10), we have

$$A_{\text{GUP}} = A + G\hbar\tilde{\alpha} \ln\left(\frac{A}{G\hbar}\right) + 4G\hbar \sum_{n=1}^{\infty} \tilde{c}_n \left(\frac{A}{G\hbar}\right)^{-n} + \text{const.} \quad (11)$$

In this case, an exact value of the constant in the right-hand side of (11) is of no great importance, as further we need the relation (1) being primarily interested in dA_{GUP} rather than in A_{GUP} ; that is, the constant in the right-hand side of (11) is insignificant. So, (1) has a fairly definite analog at high energies

$$dN_{\text{GUP}} = \frac{dA_{\text{GUP}}}{G\hbar} \quad (12)$$

that ongoing to the known low energies gives (1). There is a single considerable difference; in [14] the quantity N was defined in terms of A , and dN was defined in terms of dA , but in the case under study, the situation is opposite: A_{GUP} is defined in terms of N_{GUP} and dA_{GUP} in terms of dN_{GUP} . The logic series is here as follows:

$$A \implies N \implies N_{\text{GUP}} \implies A_{\text{GUP}}. \quad (13)$$

The GUP-correction problem of the temperature T for the arbitrary holographic screen \mathcal{S} has been studied in [70]. Actually, this case is identical to the case of the Schwarzschild black hole, and T_{GUP} was derived as a series (formulas (1), (7) from [70]):

$$T_{\text{GUP}} = T(1 + \Theta_T T^2 + \dots) = T + \Theta_T T^3 + \dots = T + \tilde{T}_{\text{GUP}}, \quad (14)$$

where the factors in the right-hand side (14) may be computed in the explicit form and at low energies $\tilde{T}_{\text{GUP}} \rightarrow 0$.

Thus, we can have a GUP - analog of Komar’s mass in ((5.33) from [14])

$$\begin{aligned} M_{\text{GUP}} &= \frac{1}{2} \int_{\mathcal{S}} T_{\text{GUP}} dN_{\text{GUP}} \\ &= \frac{1}{2} \int_{\mathcal{S}} (T + \tilde{T}_{\text{GUP}}) dN_{\text{GUP}} = \frac{1}{2G\hbar} \int_{\mathcal{S}} T_{\text{GUP}} dA_{\text{GUP}}, \end{aligned} \quad (15)$$

that the low-energy limit gives the well-known Komar formula [71], [72, page 289].

It is clear that the ‘‘GUP-deformed Komar’s mass’’ M_{GUP} in the first term (15) as a component has the known Komar’s mass [71], [72, (11.2.9–11.2.10)], [14, (5.34)]

$$M = \frac{1}{4\pi G} \int_{\mathcal{S}} T dA. \quad (16)$$

3. N_{GUP} , A_{GUP} , and M_{GUP} in Terms of Unified Small Parameter

If feasible, it is desirable to express all the above-derived fundamental quantities in terms of a unified parameter. As shown by the author in [39, 73], this is possible for black holes within the scope of GUP, and a role of the unified small parameter is played by the parameter introduced previously in [27–36] as follows:

$$\alpha_x = \frac{l_{\min}^2}{x^2}, \quad (17)$$

where x is the measuring scale, $l_{\min} \sim l_p$ by virtue of GUP (3), and $0 < \alpha \leq 1/4$.

Obviously, the principal results obtained in [39, 73] remain in force for an arbitrary screen \mathcal{S} and may be applied to the quantities N_{GUP} , A_{GUP} , and M_{GUP} defined in the preceding section.

Substituting from GUP (3) $l_{\min} = 2\sqrt{\alpha}l_p$ and using the formula $A = 4\pi R^2$, where R is the radius of the screen \mathcal{S} , we get N_{GUP} (9) of the following form:

$$N_{\text{GUP}} = N + \tilde{\alpha} \ln(\sigma\alpha_R^{-1}) + 4 \sum_{n=1}^{\infty} \tilde{c}_n \sigma^{-n} \alpha_R^n + \text{const}. \quad (18)$$

Here, α_R is a value of α parameter at the point R , and σ is equal to $16\alpha\pi$. It is convenient to refer to the form N_{GUP} derived in (18) as to the α —representation.

Using (10) and (11), we can easily obtain α —representation for A_{GUP} :

$$A_{\text{GUP}} = A + \tilde{\alpha} G \hbar \ln(\sigma\alpha_R^{-1}) + 4G\hbar \sum_{n=1}^{\infty} \tilde{c}_n \sigma^{-n} \alpha_R^n + \text{const}. \quad (19)$$

Also, it is clear that M_{GUP} (15) may be derived in terms of α_R .

Here, α_x is considered as a deformation parameter for the Heisenberg algebra ongoing from HUP to GUP. Generally speaking, initially the construction of such a deformation was realized with other parameters (e.g., [60, 61]). But it is easily shown that QFT parameter of the deformations associated with GUP may be expressed in terms of the parameter α that has been introduced in the approach to the density matrix deformation [38, 39]. Here the notation of [74] is used. Then [39, page 943],

$$[\vec{x}, \vec{p}] = i\hbar(1 + \beta^2 \vec{p}^2 + \dots), \quad (20)$$

$$\Delta x_{\min} \approx \hbar \sqrt{\beta} \sim l_p. \quad (21)$$

In this case, from (20), (21), it follows that $\beta \sim 1/\mathbf{p}^2$, and for $x_{\min} \sim l_p$, β corresponding to x_{\min} is nothing else but

$$\beta \sim \frac{1}{P_{pl}^2}, \quad (22)$$

where P_{pl} is Planck’s momentum: $P_{pl} = \hbar/l_p$.

In this manner, β is varying over the following interval:

$$\frac{\lambda}{P_{pl}^2} \leq \beta, \quad (23)$$

where λ is a numerical factor and the second term in (20) is accurately reproduced in the momentum representation (up to the numerical factor) by $\alpha_x = l_{\min}^2/x^2 \sim l_p^2/x^2 = P^2/P_{pl}^2$

$$[\vec{x}, \vec{p}] = i\hbar(1 + \beta^2 \vec{p}^2 + \dots) = i\hbar(1 + a_1 \alpha_x + a_2 \alpha_x^2 + \dots). \quad (24)$$

In the case under study, convenience of using α_x stems from its smallness, its dimensionless character, and ability to test changes in the radius R of the holographic screen \mathcal{S} .

4. Quantum Corrections to the Principal Result and Ultraviolet Limit

Based on the aforesaid, we can proceed the generalization of the results from Section 5.2 of [14] and the derivation of equations for a gravitational field within the scope of GUP.

We must consider two absolutely different cases.

4.1. Quantum Corrections to the Principal Result. It is assumed that the screen radius is given by \mathcal{S}

$$R \gg l_p. \quad (25)$$

In terms of the deformation parameter α_x introduced in the previous section, we have

$$\alpha_R \ll \frac{1}{4}. \quad (26)$$

So far we are not concerned with the redefinition of the lower limit for α_R . This is interesting when going to the infrared limit.

Then, the principal result from the final part of Section 5.2 in [14] remains valid owing to the replacement of M (formula (5.33) from [14]) by $M_{\text{GUP}} = M_{\text{GUP}}[\alpha_R]$ (15). The ‘‘ α_R -complement’’ (i.e., the difference $\bar{M}[\alpha_R] = M_{\text{GUP}}[\alpha_R] - M$) to M will be simply a (small) quantum correction for the principal result.

In this way, because in the case of GUP-corrections the left side of formula (5.33) and hence the left side in formulas (5.34), (5.35) from [14] are dependent on α_R ; the right sides of the corresponding formulas are also dependent on α_R , in particular the quantities T_{ab} , R_{ab} , and g_{ab} from (5.37).

But in fact, this relation at low energies ((25) or (26)) is insignificant since α_R at low energies (as distinct from high

Planck's energies) is varying very slowly, practically showing continuity though being discrete in character.

Indeed, as long as there is the minimal length $l_{\min} \sim l_p$, all the lengths measured are its multiples, and hence α_R is a discrete nonuniformly varying quantity. Then, due to ((25) or (26)), the difference between two successive values of α_R is as follows:

$$\Delta_{\min} [\alpha_R] = \alpha_R - \alpha_{R+l_{\min}} \sim \frac{l_{\min}^3}{R^3}, \quad (27)$$

for $R \gg l_p$ or, that is the same, for $R \gg l_{\min}$ giving a value close to zero.

And assuming in this case that α_R is continuously varying from R and all the quantities in Section 5.2 of [14] are also continuously dependent on α_R (26), we can write down the " α -analog" of formula (5.37) in [14] as

$$\begin{aligned} 2 \int_{\Sigma} \left(T_{ab} [\alpha] - \frac{1}{2} T [\alpha] g_{ab} [\alpha] \right) n^a \xi^b dV \\ = \frac{1}{4\pi G} \int_{\Sigma} R_{ab} [\alpha] n^a \xi^b dV, \end{aligned} \quad (28)$$

where the dependence of $T_{ab}[\alpha]$ and $R_{ab}[\alpha]$ on $\alpha = \alpha_R$ is completely determined, in accordance with [14, 72], by the integral $M_{\text{GUP}}[\alpha]$ (15).

Besides, it is assumed that n^a and ξ^b are dependent on α , though the dependence is dropped.

Next, similar to [14], from (28) we can derive the α -deformed Einstein Equations using the method from [1]. Note that both this method and its minor modification given ([14], end of Section 5.2) in this case are valid because α_R is small and continuous; the whole system is being continuously dependent on it.

Solutions of the α -deformed Einstein Equations represent a series in α_R , and for $\alpha_R \rightarrow 0$ or for $\alpha' = 0$ from formula (3) become the corresponding solutions of Section 5.2 in [14].

Using the result obtained in [75], we can easily extend the above result to the case with a nonzero cosmological term $\Lambda \neq 0$. In [75], Komar's formula was generalized to the case of a nonzero Λ . All the arguments from Section 5.2 of [14] in this case remain valid, and formula (5.37) takes the following form:

$$\begin{aligned} 2 \int_{\Sigma} \left(T_{ab} - \frac{1}{2} T g_{ab} \right) n^a \xi^b dV \\ = \frac{1}{4\pi G} \int_{\Sigma} (R_{ab} + \Lambda g_{ab}) n^a \xi^b dV. \end{aligned} \quad (29)$$

We can easily obtain the α -analog of the last formula with the dynamic cosmological term $\Lambda(\alpha)$ as a corresponding complement to the right-hand side (28). Analysis of the relationship between Λ and α , applicable in this case as well will be given in Section 4.2.

4.2. Ultraviolet Limit. In the case in question, we suggest that the screen \mathcal{S} has a radius on the order of several Planck's lengths

$$R \approx \xi l_{\min} = 2\alpha' \xi l_p, \quad (30)$$

where ξ is a number on the order of 1 or

$$\alpha_R \approx \frac{1}{4}. \quad (31)$$

The problem is which object puts the limit for such a screen \mathcal{S} . It may be assumed that if $T_{ab} \neq 0$ then the object may be represented only by Planck's black hole or by a microblack hole with a radius on the order of several Planck's lengths.

Clearly, the methods of [1, 14] are not in force for such screen \mathcal{S} . Specifically, there are no classical analogs of N , T , and M for the screen. Moreover, it is impossible to use the result of [1] as "a very small region the space-time" is no longer "an approximate Minkowski space" [14].

Also, such microblack hole is a horizon space, jet at high energies (Planck scales). As is known, for horizon spaces, black holes in particular, at low energies (semiclassical approximation) the results of [12] are valid.

At the horizon (and we are interested in this case only) Einstein's field Equations may be written as a thermodynamic identity ([12] formula (119)):

$$\frac{\hbar c f'(a)}{k_B T} \frac{c^3}{G \hbar} d \left(\frac{1}{4} 4\pi a^2 \right) - \frac{1}{2} \frac{c^4 da}{G} = P d \left(\frac{4\pi}{3} a^3 \right), \quad (32)$$

where $R = a$ is a radius of a black hole (i.e., of the screen \mathcal{S}), $P = T_R^R$ is the trace of the momentum-energy tensor and radial pressure, and the horizon location will be given by simple zero of the function $f(R)$, at $R = a$.

The main ingredients of (32) may be written in terms of the deformation parameter α with the coefficients containing only the numerical factors and fundamental constants [73].

Also, the work [73] presents two possible variants of high-energy (Planck) α -deformation $\alpha \rightarrow 1/4$ (32).

Hereinafter, we assume that the energy-momentum tensor of matter fields is not traceless:

$$T_a^a \neq 0, \quad (33)$$

similar, in particular, to the case under study (32) $P = T_R^R \neq 0$.

4.2.1. Case of Equilibrium Thermodynamics [73, Section 6.1].

In this case, it is assumed that in the high-energy (ultraviolet (UV)) limit the thermodynamic identity (32) is retained but now all the quantities involved in this identity become α -deformed ($\alpha \rightarrow 1/4$). All the quantities Y in (32) are replaced by the corresponding quantities Y_{GUP} with the subscript GUP. Then, the high-energy α -deformation of (32) takes the form

$$k_B T_{\text{GUP}}(\alpha) dS_{\text{GUP}}(\alpha) - dE_{\text{GUP}}(\alpha) = P(\alpha) dV_{\text{GUP}}(\alpha). \quad (34)$$

Substituting into (34) the corresponding quantities $T_{\text{GUP}}(\alpha)$, $S_{\text{GUP}}(\alpha)$, $E_{\text{GUP}}(\alpha)$, $V_{\text{GUP}}(\alpha)$, and $P(\alpha)$ and expanding them into a Laurent series in terms of α , close to high values of α , specifically close to $\alpha = 1/4$, we can derive a solution for the high energy α -deformation of the general relativity (34) as a function of $P(\alpha)$. Provided at high energies the generalization of (32) to (34) is possible; we can have the high-energy α -deformation of the metric.

It is noteworthy that in (34) T_{GUP} this time is calculated from [64, formula (10)]

$$\begin{aligned} T_{\text{GUP}}^{\text{BH}} &= \frac{1}{4\pi} \frac{\hbar R}{2\alpha'^2 l_p^2} \left[1 - \sqrt{1 - \frac{\alpha'^2 l_p^2}{R^2}} \right] \\ &= \frac{\hbar \alpha_R^{-1}}{4\pi \alpha' l_p} \left[1 - (1 - \alpha_R)^{1/2} \right] \end{aligned} \quad (35)$$

with subsequent replacement of l_p by $\sqrt{G\hbar}$ for $c = 1$.

It is especially interesting to consider the following case.

4.2.2. Case of Nonequilibrium Thermodynamics [73, Section 6.2]. In this case, the α -dependent dynamic cosmological term $\Lambda(\alpha) \neq 0$ appears in the right-hand side of (34). Then, with the addition of $\Lambda(\alpha) \neq 0$, the α -representation (34) (for $\hbar = 1$) is given as follows [73, formula (53)]:

$$-\alpha^2 f'(\alpha) - \frac{1}{2}\alpha = 16\pi\alpha'^2 P(\alpha) G^2 - G\Lambda(\alpha), \quad (36)$$

where $\alpha = \alpha_R \approx 1/4$,

$$f'(\alpha) = 4\pi k_B T_{\text{GUP}}(\alpha), \quad (37)$$

and the derivative in the left-hand side of (37) is taken with respect to α .

$\Lambda(\alpha)$ in the right-hand side of (37) may be subjected to a series expansion in terms of α , in compliance with the holographic principle [49–53] as applied to the Universe [76]. In [38, 39, 73, 77] in the leading order, this expansion results in the first power, that is, we have

$$\Lambda(\alpha_R) \sim \alpha_R \Lambda_p, \quad (38)$$

where Λ_p is the initial value of $\Lambda \approx \Lambda_{1/4}$ derived using the well-known procedure of “summation over all zero modes” and the Planck momentum cutoff [78, 79]. Actually, (38) is in a good agreement with the observable $\Lambda = \Lambda_{\text{observ}}$. Because a radius of the visible part of the Universe is given as $R = R_{\text{Univ}} \approx 10^{28}$ cm, it is clear that $\alpha_R \approx 10^{-122}$, and (38) is completely consistent with the experiment [79].

Note that proceeding directly from a quantum field theory but without the use of the holographic principle, we can have only a rough estimate of Λ that, on the whole, is at variance with Λ_{observ} . Such an estimate may be obtained in different ways: by simulation [80]; using the cutoff [78] but now in the infrared limit; with the use of the Generalized Uncertainty Principle for the pair (Λ, V) , where V is four-dimensional volume [38, 39]. In the α -representation in this case, the expansion in terms of α results in the second leading order

$$\Lambda(\alpha_R) \sim \alpha_R^2 \Lambda_p, \quad (39)$$

that, obviously, is at variance with the accepted facts.

5. Conclusion

(I) A very interesting case of the zero energy-momentum tensor for matter fields $T_{ab} = 0$ and, specifically the case of $P(\alpha) = 0$ in the right-hand side of (36) has remained beyond the scope of the final section. We can state the problem more specifically: for which conditions in this case we can derive a solution in the form of the de Sitter space with large values of α ?

This problem is also important when we try to find whether it is possible to derive the initial inflation conditions [81, 82] for $T_{ab} = 0$ on the basis of the foregoing analysis.

Note that the dynamic cosmological term $\Lambda(\alpha)$ correlates well with inflation models [81, 82] as the latter require a very high Λ at the early stages of the Universe, and this is distinct from $\Lambda = \Lambda_{\text{exper}}$ in the modern period. Of great interest is the recent work [83], where a mechanism of the vacuum energy decay in the de Sitter space is established to support a dynamic nature of Λ .

(II) The deformation parameter α_R has a double meaning.

As $\kappa = 1/R$ —curvature with the radius R , $\alpha_R = \kappa^2 l_{\text{min}}^2$ is nothing else but the squared curvature multiplied by the squared minimal area and is explicitly dependent on the energy E . On the other hand, it is seen that, at least at the known energies, from the definition of the bit number N in [14, formula (3.10)], we get $\alpha_R \sim 1/N$. But, because $\alpha_R = \alpha_R[E]$, this suggests that $N = N[E]$, as demonstrated in the text ongoing to higher energies

$$N \implies N_{\text{GUP}}. \quad (40)$$

Nevertheless, ongoing to lower energies, that is, in the infrared limit, the same should be true: the bit number must be a function of energy.

This problem and the relevant questions touched upon in this paper will be further considered in subsequent works of the author.

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