On strong coupling in nonrelativistic general covariant theory of gravity

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We study the strong coupling problem in the Horava-Melby-Thompson setup of the Horava-Lifshitz gravity with an arbitrary coupling constant $\lambda$, generalized recently by da Silva, where $\lambda$ describes the deviation of the theory in the infrared from general relativity that has $\lambda_{GR} = 1$. We find that a scalar field in the Minkowski background becomes strong coupling for processes with energy higher than $\Lambda_{\infty} \equiv (M_{pl}/c_1)^{3/2}M_{pl}(\lambda - 1)^{5/4}$, where generically $c_1 < M_{pl}$. However, this problem can be cured by introducing a new energy scale $M$, so that $M < \Lambda_{\infty}$, where $M$ denotes the suppression energy of high order derivative terms of the theory.

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I. INTRODUCTION

Horava recently proposed a new theory of quantum gravity \[1\], based on the perspective that Lorentz symmetry may appear as an emergent symmetry at low energies, but can be fundamentally absent at high energies. His starting point is the anisotropic scalings of space and time,

$$x \to b^{-1}x, \quad t \to b^{-2}t.$$  \hspace{1cm} (1.1)

In $(3+1)$-dimensions, in order for the theory to be power-counting renormalizable, the critical exponent $z$ needs to be $z \geq 3$ \[2\]. At long distances, all the high-order curvature terms are negligible, and the linear term $R$ becomes dominant. Then, the theory is expected to flow to the relativistic fixed point $z = 1$, whereby the general covariance is “accidentally restored.” The special role of time is realized with the Arnowitt-Deser-Misner decomposition \[3\],

$$ds^2 = -N^2 \dot{t}^2 + g_{ij} (dx^i + N^i dt) (dx^j + N^j dt),$$  \hspace{1cm} (1.2)

Under the rescaling \[1,1\] with $z = 3$, $N$, $N^i$ and $g_{ij}$ scale, respectively, as,

$$N \to N, \quad N^i \to b^{-2} N^i, \quad g_{ij} \to g_{ij}.$$  \hspace{1cm} (1.3)

The gauge symmetry of the system now is broken down to the foliation-preserving diffeomorphisms $\text{Diff}(M, F)$,

$$\delta t = -f(t), \quad \delta x^i = -\dot{\zeta}^i(t, x),$$  \hspace{1cm} (1.4)

under which, $N$, $N^i$ and $g_{ij}$ transform as,

$$\delta N = \zeta^k \nabla_k N + \dot{N} f + \dot{f},$$

$$\delta g_{ij} = \nabla_i \dot{\zeta}^j + \nabla_j \dot{\zeta}^i + f g_{ij},$$  \hspace{1cm} (1.5)

where $f \equiv df/dt$, $\nabla_i$ denotes the covariant derivative with respect to the 3-metric $g_{ij}$, $\dot{N} = g_{ik} \dot{\zeta}^k$, and $\delta g_{ij} = g_{ij}(t, x^k) - g_{ij}(t, x^k)$, etc. Eq. (1.5) shows clearly that the lapse function $N$ and the shift vector $\dot{N}$ play the role of gauge fields of the $\text{Diff}(M, F)$. Therefore, it is natural to assume that $N$ and $N^i$ inherit the same dependence on space and time as the corresponding generators,

$$N = N(t), \quad N^i = N_i(t, x),$$  \hspace{1cm} (1.6)

while the dynamical variables $g_{ij}$ in general depend on both time and space, $g_{ij} = g_{ij}(t, x)$. This is often referred to as the projectability condition.

Abandoning the general covariance, on the other hand, gives rise to a proliferation of independently coupling constants, which could potentially limit the prediction powers of the theory. Inspired by condensed matter systems \[4\], Horava assumed that the gravitational potential $\mathcal{L}_V$ can be obtained from a superpotential $W_g$ via the relations,

$$\mathcal{L}_{V, \text{detailed}} = w^2 E_{ijkl} \mathcal{G}^{ijkl} E_{kl},$$

$$E^{ij} = \frac{1}{\sqrt{g}} \frac{\delta W_g}{\delta g_{ij}},$$  \hspace{1cm} (1.7)

where $w$ is a coupling constant, and $\mathcal{G}^{ijkl}$ denotes the generalized De Witt metric, defined as $\mathcal{G}^{ijkl} = (g^{ik} g^{jl} + g^{il} g^{jk})/2 - \lambda g^{ij} g^{kl}$, with $\lambda$ being a coupling constant. The general covariance, $\delta x^\mu = \zeta^\mu(t, x)$, $(\mu = 0, 1, ..., 3)$, requires $\lambda = 1$. The superpotential $W_g$ is given by

$$W_g = \int_{\Sigma} \omega_3(\Gamma) + \frac{1}{\kappa_W^2} \int d^3 x \sqrt{g}(R - 2\Lambda),$$  \hspace{1cm} (1.8)

with $\omega_3(\Gamma)$ being the gravitational Chern-Simons term,

$$\omega_3(\Gamma) = \text{Tr} \left( \Gamma \wedge d\Gamma + \frac{2}{3} \Gamma \wedge \Gamma \wedge \Gamma \right).$$  \hspace{1cm} (1.9)
The condition (1.7) is usually referred to as the detailed balance condition.

However, with this condition it was found that the Newtonian limit does not exist [8], and a scalar field in the UV is not stable [9]. Thus, it is generally believed that this condition should be abandoned [10]. But, it has several remarkable features [8]: it is in the same spirit of the AdS/CFT correspondence [9]; and in the nonequilibrium thermodynamics, the counterpart of the superpotential $W_0$ plays the role of entropy, while the term $E^{ij}$ the entropic forces [10]. This might shed light on the nature of the gravitational forces, as proposed recently by Verlinde [11]. Due to these desired properties, together with Borzou, two of the present authors recently studied this condition in detail, and found that the scalar field can be stabilized, if the detailed balance condition is allowed to be softly broken [12]. This can also solve the other problems [5, 7]. In addition, such a breaking is allowed to be softly broken [12]. This can also solve the instability, ghost, and strong coupling [13–17]. To overcome those problems, recently Horava and Melby-Thompson (HMT) [18] extended the foliation-preserving-diffeomorphisms $\text{Diff}(M, \mathcal{L})$ to include a local $U(1)$ symmetry.

$$U(1) \ltimes \text{Diff}(M, \mathcal{L}).$$

(1.10)

Such an extended symmetry is realized by introducing a $U(1)$ gauge field $A$ and a Newtonian prepotential $\varphi$. Under $\text{Diff}(M, \mathcal{L})$, these fields transform as $\text{Diff}(\mathcal{M}, \mathcal{F})$.

$$\delta A = \zeta^i \partial_i A + f A + f A, \quad \delta \varphi = f \dot{\varphi} + \zeta^i \partial_i \varphi,$$

(1.11)

while under $U(1)$, characterized by the generator $\alpha$, they, together with $N$, $N^i$ and $g_{ij}$, transform as

$$\delta_\alpha A = \alpha - N^i \nabla_i \alpha, \quad \delta_\alpha \varphi = -\alpha, \quad \delta_\alpha N^i = N \nabla_i \alpha, \quad \delta_\alpha g_{ij} = 0 = \delta_\alpha N.$$

(1.12)

HMT showed that, similar to GR, the spin-0 graviton is eliminated [18]. This was further confirmed in [19]. Then, the instability of the spin-0 gravity is out of question. In addition, in the linearized theory Horava noticed that the $U(1)$ symmetry only pertains to the case $\lambda = 1$ [1]. It was believed that this was also the case when the Newtonian prepotential is introduced [18]. If this were true, the ghost and strong coupling problems would be also resolved, as both of them are related to the very fact that $\lambda \neq 1$ [20].

However, it has been soon challenged by da Silva [21], who argued that the introduction of the Newtonian prepotential is so strong that actions with $\lambda \neq 1$ also has the extended symmetry, Eq. (1.10). Although the spin-0 graviton is also eliminated even with any $\lambda$ as shown explicitly in [20, 22], the ghost and strong coupling problems rise again, because now $\lambda$ can be different from one. Indeed, it was shown [20] that to avoid the ghost problem, $\lambda$ must satisfy the same constraints,

$$\lambda \geq 1, \quad \text{or} \quad \lambda < \frac{1}{3}.$$

(1.13)

as found previously [1, 23–25].

In addition, the strong coupling problem also arise [20]. In this paper, we shall address this important issue. In particular, in Sec. II we briefly review the HMT setup with any coupling constant $\lambda$ with detailed balance condition softly breaking, a version presented in [12] in detail, while in Sec. III we study the strong coupling problem when a scalar field is present. We find that a scalar field in the Minkowski background becomes strong coupling for processes with energy higher than $\Lambda_\varphi \equiv (M_{pl} / c_1)^{1/2} M_{pl} |\lambda - 1|^{1/4}$. For $c_1 \approx M_{pl}$, this gives precisely the strong coupling strength found in [20]. However, this problem can be resolved by introducing a new energy scale $M_s [20]$, so that $M_s \lesssim \Lambda_\varphi$, where $M_s$ defines the energy scale that suppresses the sixth-order derivative terms of the theory. In Sec. IV, we present our main conclusions.

It should be noted that the strong coupling problem in other versions of the Horava-Lifshitz (HL) theory has been studied extensively by using both effective field theory [20, 27–30] and Stückelberg formalism [26, 31, 32]. In this paper, we shall follow the approach of the effective field theory [33], although the final conclusions are independent of the methods to be used. In addition, strong coupling can happen not only due to gravitational/matter self-interactions, but also to the interactions among gravitational and matter fields. The latter was studied in [20, 30], while strong coupling due to the self-gravitational interactions were studied in [20, 24, 32]. In this paper, we shall study it due to the interaction between gravitational and scalar fields.

II. GENERAL COVARIANT THEORY WITH DETAILED BALANCE CONDITION SOFTLY BREAKING

As mentioned above, HMT considered only the case $\lambda = 1$. Later, da Silva generalized it to the cases with any $\lambda$ [21], in which the total action can be written in the form [20, 21],

$$S = \zeta^2 \int dt \sqrt{g} \left( \mathcal{L}_K - \mathcal{L}_N - \mathcal{L}_\varphi + \mathcal{L}_A + \mathcal{L}_\lambda + \zeta^2 \mathcal{L}_M \right),$$

(2.1)

where $g = \det g_{ij}$,

$$\mathcal{L}_K = K_{ij} K^{ij} - \lambda K^2,$$

$$\mathcal{L}_\varphi = \varphi G^{ij} \left( 2 K_{ij} + \nabla_i \nabla_j \varphi \right),$$
\[ L_A = \frac{A}{N} (2\Lambda_g - R), \]
\[ L_\lambda = (1 - \lambda) \left[ (\Delta \varphi)^2 + 2K\Delta \varphi \right], \]
\[ \Delta \equiv \nabla^2 = g^{ij} \nabla_i \nabla_j, \]
and \( \Lambda_g \) is a coupling constant. The Ricci and Riemann terms all refer to the 3-metric \( g_{ij} \), and
\[ K_{ij} = \frac{1}{2N} (\delta_{ij} + \nabla_i N_j + \nabla_j N_i), \]
\[ G_{ij} = R_{ij} - \frac{1}{2} g_{ij} R + \Lambda_g g_{ij}. \]

\( L_V \) is an arbitrary Diff(\( \Sigma \))-invariant local scalar functional built out of the spatial metric, its Riemann tensor and spatial covariant derivatives. With the detailed balance condition softly breaking, it takes the form 12, 34,
\[ L_V = \zeta^2 \gamma_0 + \gamma_1 R + \frac{1}{\xi^3} \left( \gamma_2 R^2 + \gamma_3 R_{ij} R^{ij} \right) \]
\[ + \frac{\gamma_4}{\xi^4} \epsilon^{ijk} R_{ij} \nabla_k R_{ij} + \frac{\gamma_5}{\xi^4} C_{ij} C^{ij}, \]
where the coupling constants \( \gamma_s (s = 0, 1, 2, \ldots 5) \) are all dimensionless, and \( \gamma_5 \equiv w^2 \xi^4 \). The relativistic limit in the IR requires
\[ \lambda = 1, \quad \gamma_1 = -1, \quad \zeta^2 = \frac{1}{16\pi G}. \]

The existence of the \( \gamma_4 \) term explicitly breaks the parity, which could have important observational consequences on primordial gravitational waves 23.

\( L_M \) is the matter Lagrangian density. For a scalar field \( \chi \) with detailed balance conditions softly breaking, it is given by 12,
\[ L_M = L_\chi^{(A, \varphi)} + L_\chi^{(0)}, \]
where
\[ L_\chi^{(A, \varphi)} = \frac{A - A}{N} \left[ c_1 \Delta \chi + c_2 (\nabla \chi)^2 \right] \]
\[ + \frac{1}{N} \left( \chi - N^i \nabla_i \chi \right) \left( \nabla^k \chi \right) \nabla_k \chi \]
\[ + \frac{f}{2} \left( (\nabla^k \chi) (\nabla_k \chi) \right)^2, \]
\[ L_\chi^{(0)} = \frac{f}{N^2} \left( \chi - N^i \nabla_i \chi \right)^2 - V, \]
and
\[ V = V(\chi) + \frac{1}{2} V_i (\chi) (\nabla \chi)^2 + V_2 (\chi) P_i^2 \]
\[ + V_3 (\chi) P_i^3 + V_4 (\chi) P_i + V_5 (\chi) (\nabla \chi)^2 P_i \]
\[ + V_6 P_1 P_2, \]
with \( V(\chi) \) and \( V_n(\chi) \) being arbitrary functions of \( \chi \), and
\[ P_n \equiv \Delta^n \chi, \quad V_6 \equiv -\sigma_3^2. \]

The corresponding field equations are given in Appendix A and by Eqs.(3.13) and (3.14) in 12.

### III. STRONG COUPING

To study the strong coupling problem, it is found sufficient to consider perturbations of the Minkowski spacetime,
\[ \bar{N} = 1, \quad \bar{N}_i = 0, \quad \bar{g}_{ij} = \delta_{ij}, \quad \bar{A} = \bar{\varphi} = 0, \quad \bar{\chi} = \bar{\chi}_0, \quad V(\bar{\chi}_0) = V'(\bar{\chi}_0) = 0 = \Lambda = \Lambda_g, \]
where \( \chi_0 \) is a constant. Without loss of generality, we can set \( \chi_0 = 0 \), a condition that will be assumed in the rest of the paper. The perturbed fields with the (generalized) quasilongitudinal gauge are given by 19, 20, 25,
\[ N = 1, \quad N_i = B_i, \quad g_{ij} = (1 - 2\psi) \delta_{ij}, \quad A = \delta A, \quad \varphi = 0, \quad \chi = \delta \chi. \]

Then, we find
\[ \sqrt{g} L_K = (1 + \psi) \left[ (1 - 3\lambda)(3\psi^2 + 2\psi \partial^2 B) + B_{ij} B^{ij} \right] \]
\[ - \lambda (\partial^2 B)^2 + 2 B_{ij} B^{ij} \psi^2 - (1 - 3\lambda) \psi B_i \psi^i - (1 - \lambda) B_{ij} \psi^i \psi^j, \]
\[ \sqrt{g} L_A = -2A \left[ 2(1 + \psi) \partial^2 \psi + 3\psi, \psi^i \right], \]
\[ \sqrt{g} L_V = 4\partial^2 \psi + (2) \left( \sqrt{g} L_\chi \right) + (3) \left( \sqrt{g} L_\chi \right), \]
where the first order term \( \partial^2 \psi \) in the expression of \( \sqrt{g} L_\chi \) becomes a boundary term after integration, and therefore can be discarded, and
\[ (2) \left( \sqrt{g} L_\chi \right) = 2 \psi \partial^2 \psi + \frac{2}{\xi^2} (8\gamma_2 + 3\gamma_3) \psi \partial^4 \psi, \]
\[ (2) \left( \sqrt{g} L_\chi \right) = \frac{f}{2} \chi^2 - \frac{1}{2} \sqrt{V''} \chi^2 + \frac{1}{2} (1 + 2V_1) \chi \partial^2 \chi \]
\[ - (V_2 + V_1') \chi \partial^4 \chi + \sigma_3^2 \chi \partial^6 \chi \]
\[ + c_1 \chi \partial^2 A, \]
\[ (3) \left( \sqrt{g} L_\chi \right) = -6(\psi^2 \partial^2 \psi + 3\psi \psi, \psi^i) \]
\[ + A_1 (\gamma_2, \gamma_3) \psi^2 \partial^4 \psi + O(\gamma_2, \gamma_3, \psi, \partial^i), \]
\[ + A_2 (\gamma_5) \psi^2 \partial^6 \psi + O(\gamma_5, \psi, \partial^6), \]
\[ (3) \left( \sqrt{g} L_\chi \right) = -\frac{3}{2} f \psi \chi^2 - f \chi B_i \chi^i - c_1 \psi A \partial^2 \chi \]
\[ - c_1 A \psi^i \chi^i \chi^i + c_2 A \chi^i \chi^i \chi^i \chi^i + \frac{3}{2} V'' \psi \chi^2. \]
where the terms of $A_i$ ($B_i$) are representatives of the fourth and sixth order derivative terms of $\psi$ ($\chi$), and the specific dependences of them on their arguments are not important to the analysis of the strong coupling problem, as shown below. So, for the sake of simplicity, we shall not give them explicitly here. Hence, to second order we obtain

$$S^{(2)} = \zeta^2 \int dt d^3x \left\{ 3(1 - 3\lambda)\dot{\psi}^2 + 2(1 - 3\lambda)\dddot{\psi}\ddot{B} + (1 - \lambda)(\ddot{\psi}\dot{B})^2 - 4\dddot{\psi}\ddot{B} + 2\dddot{\psi} - \frac{2}{\zeta^2}(8\gamma_2 + 3\gamma_3)\dot{\psi}\dot{\chi} + \frac{1}{\zeta^2} \left\{ \frac{1}{2} f \dot{\chi}^2 - \frac{1}{2} V'' \dot{\chi}^2 + \frac{1}{2}(1 + 2V_1)\chi\ddot{\chi} - (V_2 + V_3')\chi\dot{\chi}^2 + \sigma_3^2 \chi\dot{\chi}^2 \right\} \right\},$$  

(3.7)

Variations of $S^{(2)}$, respectively, with respect to $\psi$, $B$, $A$ and $\chi$ yield,

$$\dddot{\psi} + \frac{1}{3}\dot{\psi}\dot{B} = \frac{2}{3(3\lambda - 1)} \partial^2 \left( A + \psi \right) + \frac{8\gamma_2 + 3\gamma_3}{\zeta^2} \partial^2 \psi,$$

(3.8)

$$(1 - 3\lambda)\psi + (1 - \lambda)\dddot{B} = 0,$$

(3.9)

$$\psi = 4\pi Gc_1 \chi,$$

(3.10)

and

$$f \dddot{\chi} + V'' \dddot{\chi} - (1 + 2V_1)\dddot{\chi} + 2(V_2 + V_3')\dddot{\chi} - 2\sigma_3^2 \dddot{\ chi} = c_1 \dddot{A}. $$

(3.11)

The above equations can be obtained from Eqs.(4.13) - (4.17) and Eq.(4.20) of [12], by setting $\dddot{x} = \dddot{\chi} = 0$ and $\alpha = 1$, as it is expected. Using Eqs.(3.8)-(3.9), we can integrate out $\psi$, $B$ and $A$, so $S^{(2)}$ finally takes the form,

$$S^{(2)} = \beta^2 \int dt d^3x \left\{ \dot{\chi}^2 - \frac{\alpha}{\beta} (\partial \chi)^2 - m_\chi^2 \chi^2 \right\} - \frac{1}{M_A} \chi\partial^4 \chi + \frac{1}{M_B} \chi\partial^6 \chi,$$

(3.12)

where

$$\beta^2 = \frac{2\pi Gc_1^2}{|c_\psi|^2} + \frac{f}{2},$$

$$\alpha = \frac{1}{2\pi^2}(1 + 2V_1 - 4\pi Gc_1^2),$$

$$M_A^\dagger = \beta^2 \left( \frac{2\pi Gc_1^2 8\gamma_2 + 3\gamma_3}{\zeta^2} + V_2 + V_4' \right)^{-1},$$

$$M_B^\dagger = \frac{\beta^2 c_\psi^2}{\zeta^2},$$

and $c_\psi^2 = (1 - \lambda)/(3\lambda - 1)$. As a consistency check, one can show that the variation of the action (3.12) with respect to $\chi$ yields the master equation (4.21) given in [12]. In addition, when $\lambda$ satisfies the condition (1.13), the above expression shows clearly that the scalar field is ghost free for $f > 0$, as first noticed in [12]. The scalar field is stable in all energy scales by properly choosing the potential terms $V_i$, including the UV and IR. For detail, we refer readers to [12].

Therefore, in the following we focus only on the strong coupling problem. To this end, let us first note that the cubic action is given by $^1$,

$$S^{(3)} = \int dt d^3x \left\{ \lambda_1 \left( \frac{1}{\partial^2 \dot{\chi}} \right) \chi\partial^2 \chi + \lambda_2 \left( \frac{1}{\partial^2 \dot{\chi}} \right) \partial^4 \chi \right\},$$

(3.14)

where $\ldots$ represents the fourth- and sixth-order derivative terms of $O$'s given in Eq.(3.6), which are irrelevant to the strong coupling problem, as mentioned above, and

$$\lambda_1 = \frac{c_1^3}{8\xi^3|c_\psi|^2|c_\psi|^2}, \quad \lambda_2 = \frac{1}{|c_\psi|^2} \left( \frac{5c_1^3}{32c_\psi^4} - \frac{c_1c_2}{4c_\psi^3} \right),$$

$$\lambda_3 = \frac{c_1^3}{32\xi^3|c_\psi|^2 - \frac{3c_1}{4c_\psi}, \quad \lambda_4 = \frac{c_1^3}{64\xi^3|c_\psi|^2},$$

$$\lambda_5 = \frac{c_1^3}{64\xi^3|c_\psi|^2 + \frac{c_1}{4c_\psi^2|c_\psi|^2}}, \quad \lambda_6 = \frac{3c_1}{8\xi^3|c_\psi|^2} - \frac{\dot{\psi}}{6},$$

$^1$ Note the difference between the perturbations considered in [12, 24, 28, 29] and the ones studied here. In particular, in [19, 20, 28, 29] the perturbations of the form, $N = 1$, $N_i = \partial_i \beta, q_{ij} = e^{2\beta}g_{ij}$, were studied, while in this paper we use the expansions of Eq.(3.3) to calculate the third order action. Although such obtained $S^{(3)}$ is different from the one obtained in [12, 20, 28, 29], it is not difficult to argue that the final conclusion for the strong coupling problem will be the same.
\[
\lambda_7 = \frac{\dot{V}_1}{2} + \frac{c_1^2 c_2}{8 \xi^2} - \frac{c_1}{16 \xi^2} \frac{c_1 V_1}{8 \xi^2}, \quad \lambda_8 = \tilde{A}_1(\gamma_2, \gamma_3, c_1) + B_1(V_2, V_4), \quad \lambda_9 = \tilde{A}_2(\gamma_5, c_1) + B_2(V_3, V_5, V_6),
\]

where \( \tilde{A}_i = (4\pi G G_0)^3 A_i \). Depending on the energy scales, each term in Eq. (3.14) will have different scalings. Thus, in the following we consider them separately.

**A. \(|\nabla| \ll M_*\)**

When \(|\nabla| \ll M_*\), where \( M_* = \text{Min}(M_A, M_B) \), we find that the high order derivative terms in Eq. (3.12) can be neglected, and

\[
S^{(2)} \simeq \beta^2 \int dt d^3 x \left[ \chi^2 - a (\partial \chi)^2 \right].
\]

(3.16)

Note that in writing the above expression, without loss of generality, we had assumed that \(|\nabla| \gg m_\chi\). Setting

\[
t = \tilde{b}_1 \tilde{t}, \quad x^i = \tilde{b}_2 \tilde{x}^i, \quad \chi = \tilde{b}_3 \hat{\chi},
\]

(3.17)

we can bring Eq. (3.16) into its “canonical” form,

\[
S^{(2)} \simeq \int dt d^3 \hat{x} \left[ \left( \partial \hat{\chi}^* \right)^2 - \left( \partial \hat{\chi} \right)^2 \right],
\]

(3.18)

in which the coefficient of each term is order of one, for

\[
b_2 = \tilde{b}_1 \sqrt{a}, \quad b_3 = \frac{1}{\tilde{b}_1 \beta^3 a^{3/4}},
\]

(3.19)

where \( \hat{\chi}^* \equiv d\hat{\chi}/d\hat{t} \). Note that the requirement that the coefficient of each term be order of one is important in order to obtain a correct coupling strength \([20, 26, 27]\). In addition, the transformations (3.17) should not be confused with the gauge choice (3.2), as it just provides a technique to obtain the correct coupling strength. In fact, when we consider physics, we will all refer to the ones obtained in the \( t \) and \( x \) coordinates, as to be shown below. Inserting Eq. (3.17) into Eq. (3.14), we obtain

\[
S^{(3)} = \frac{1}{b_1 \beta^3 a^{3/4}} \hat{S}^{(3)},
\]

(3.20)

where

\[
\hat{S}^{(3)} = \int dt d^3 \hat{x} \left( \lambda_1 \left( \frac{1}{\partial^2} \hat{\chi} \right)^* \right) \hat{\chi} \partial^2 \hat{\chi} + \lambda_2 \left( \frac{1}{\partial^2} \hat{\chi} \right)^* \partial_\tilde{t} \hat{\chi} \partial^2 \hat{\chi} + \lambda_3 \hat{\chi} \partial^2 \hat{\chi}^2 \right) + \lambda_4 \left( \partial_\tilde{t} \hat{\chi} \right) \left( \frac{\partial \hat{\chi}}{\partial^2} \right)^* \left( \frac{\partial^2 \hat{\chi}}{\partial^2} \hat{\chi}^* \right) + \lambda_5 \hat{\chi} \left( \frac{\partial \hat{\chi}}{\partial^2} \hat{\chi}^* \right) \partial_\tilde{t} \hat{\chi} + \lambda_6 \hat{\chi}^2 \partial^2 \hat{\chi}^2 + \lambda_7 \hat{\chi}^3 \partial^2 \hat{\chi}^3 + \lambda_8 \hat{\chi}^4 \partial^2 \hat{\chi}^4 + \lambda_9 \hat{\chi}^5 \partial^2 \hat{\chi}^5 \right).
\]

(3.21)

On the other hand, from Eq. (3.18), one finds that \( S^{(3)} \) is invariant under the rescaling,

\[
t \to b^{-1} \tilde{t}, \quad x^i \to b^{-1} \tilde{x}^i, \quad \hat{\chi} \to b \hat{\chi},
\]

(3.22)

while the terms of \( \lambda_1, \ldots, 5 \) and \( \lambda_7 \) in \( S^{(3)} \) all scale as \( b \), and the terms of \( \lambda_6 \) scale as \( b^{-1} \), \( b^3 \), \( b^5 \), respectively. Therefore, except for the \( \lambda_9 \) term, all the others are irrelevant and nonrenormalizable \([32]\). For example, considering a process with an energy \( E \), then we find that the fourth term has the contribution,

\[
\int dt d^3 \hat{x} \left( \partial_\tilde{t} \hat{\chi} \right) \left( \partial_\tilde{t} \left( \frac{\partial \hat{\chi}}{\partial^2} \hat{\chi}^* \right) \right) \left( \frac{\partial \hat{\chi}}{\partial^2} \hat{\chi}^* \right) \simeq E.
\]

(3.23)

Since the action \( S^{(3)} \) is dimensionless, we must have

\[
\frac{\lambda_4}{b_1 \beta^3 a^{3/4}} \int dt d^3 \hat{x} \left( \partial_\tilde{t} \hat{\chi} \right) \left( \partial_\tilde{t} \left( \frac{\partial \hat{\chi}}{\partial^2} \hat{\chi}^* \right) \right) \left( \frac{\partial \hat{\chi}}{\partial^2} \hat{\chi}^* \right) \simeq \frac{E}{\Lambda^{(4)}_{SC}},
\]

(3.24)

where \( \Lambda^{(4)}_{SC} \) has the same dimension of \( E \), and is given by,

\[
\Lambda^{(4)}_{SC} = \frac{b_1 \beta^3 a^{3/4}}{\lambda_4}.
\]

(3.25)

Similarly, one can find \( \Lambda^{(n)}_{SC} \) for all the other nonrenormalizable terms. But, when \( \lambda \to 1 \) (or \( \kappa \to 0 \)), the lowest one of the \( \Lambda^{(n)}_{SC} \)'s is given by \( \Lambda^{(4)}_{SC} \), so we have

\[
\Lambda_{\omega} \equiv \frac{b_1 \beta^3 a^{3/4}}{\lambda_4},
\]

(3.26)

above which the nonrenormalizable \( \lambda_4 \) term becomes larger than unit, and the process runs into the strong coupling regime. Back to the physical coordinates \( t \) and \( x \), the corresponding energy and momentum scales are given, respectively, by

\[
\Lambda_{\omega} = \frac{\Lambda_{\omega}}{b_1} \simeq O(1) \left( \frac{c_1}{c_5} \right)^{3/2} M_{pl} |c_\omega|^5/2, \quad \Lambda_{k} = \frac{\Lambda_{k}}{b_2} \simeq O(1) \left( \frac{c_1}{c_5} \right)^{1/2} M_{pl} |c_\omega|^{3/2}. \]

(3.27)

In particular, for \( c_1 \simeq \zeta \), we find that \( \Lambda_{\omega} \simeq M_{pl} |c_\omega|^5/2 \), which is precisely the result obtained in \([20]\).
To study the strong coupling problem, we shall follow what we did in the last case, by first writing $S^{(2)}$ in its canonical form,

$$S^{(2)} = \int dt d^{3}x \left( \tilde{\chi}^{* 2} - \tilde{\chi} \tilde{\partial}^{4} \tilde{\chi} \right),$$

through the transformations (3.17). It can be shown that now $b_{2}$ and $b_{3}$ are given by

$$b_{2} = \sqrt{\frac{b_{1}}{M_{A}}}, \quad b_{3} = \frac{M_{A}^{3/4}}{b_{1}^{1/4} \beta},$$

for which the cubic action $S^{(3)}$ takes the form,

$$S^{(3)} = \frac{M_{A}^{3/4}}{b_{1}^{1/4} \beta^{3}} \tilde{S}^{(3)},$$

where $\tilde{S}^{(3)}$ is given by Eq. (3.21). Due to the nonrelativistic nature of the action (3.30), its scaling becomes anisotropic,

$$\hat{t} \to b^{-2} \hat{t}, \quad \hat{x}^{i} \to b^{-1} \hat{x}^{i}, \quad \hat{\chi} \to b^{1/2} \hat{\chi}.$$ (3.33)

Then, we find that the first five terms in Eqs. (3.32) and (3.21) scale as $b^{1/2}$, while the terms of $\lambda_{6}, ..., \beta$ scale, respectively, as $b^{-7/2}, b^{-3/2}, b^{1/2}, b^{5/2}$. Thus, except for the $\lambda_{6}$ and $\lambda_{7}$ terms, all the others are not renormalizable. It can be also shown that the processes with energy higher than $\Lambda_{\omega}^{(A)}$ become strong coupling, where $\Lambda_{\omega}^{(A)}$ is given by

$$\Lambda_{\omega}^{(A)} = \frac{16}{81} \left( \frac{M_{pl}}{M_{A}} \right)^{3} M_{pl} |\psi|^{4}, \quad (M_{A} < M_{B}).$$ (3.34)

Therefore, when the fourth-order derivative terms dominate, the strong coupling problem still exists. This is expected, as power-counting tells us that the theory is renormalizable only when $z \geq 3$ [cf. Eq. (3.15)]. Indeed, as to be shown below, when the sixth order derivative terms dominate, the strong coupling problem does not exist any longer.

2. $M_{A} \gtrsim M_{B}$

In this case, we have $M_{A} = M_{B}$, and for processes with $E \gtrsim M_{B}$, Eq. (3.12) reduces to

$$S^{(2)} = \beta^{2} \int dt d^{3}x \left( \tilde{\chi}^{* 2} - \frac{1}{M_{B}^{2}} \chi \tilde{\partial}^{2} \chi \right).$$ (3.35)

Then, by the transformations (3.17) with

$$b_{2} = \frac{b_{1}^{1/3}}{M_{B}^{2/3}}, \quad b_{3} = \frac{M_{B}}{\beta},$$ (3.36)

we obtain,

$$S^{(2)} = \int dt d^{3}x \left( \tilde{\chi}^{* 2} - \tilde{\chi} \tilde{\partial}^{2} \tilde{\chi} \right),$$ (3.37)
while the cubic action $S^{(3)}$ becomes,

$$ S^{(3)} = \frac{M_B}{\beta^3} S^{(3)}. $$  \hspace{1cm} (3.38) 

Eq. (3.37) is invariant under the rescaling,

$$ t \to b^{-3} \hat{t}, \quad \hat{x} \to b^{-1} \hat{x}, \quad \hat{\chi} \to \hat{\chi}. $$  \hspace{1cm} (3.39) 

Then, it can be shown that the first five terms in Eqs. (3.38) and (3.21) are scaling-invariant, and so the last term. The terms of $\lambda_6, \lambda_7, \lambda_8$, on the other hand, scale, respectively, as $b^{-6}, b^{-4}, b^{-2}$. Therefore, the first five and the last terms now all become strictly renormalizable, while the $\lambda_6, \lambda_7$ and $\lambda_8$ terms become superrenormalizable [33]. To have these strictly renormalizable terms to be weakly coupling, we require their coefficients be less than unit,

$$ \frac{M_*}{\beta^3} \lambda_n < 1, \quad (n = 1, \ldots, 5, 9). $$  \hspace{1cm} (3.40) 

For $\lambda \sim 1$ (or $|c_\psi| \sim 0$), we find that the above condition holds for

$$ M_* < \frac{2}{3} M_{pl} |c_\psi|. $$

It can be shown that this condition holds identically, provided that $M_* < \Lambda_\omega$, that is,

$$ M_* < \left( \frac{\varsigma}{c_1} \right)^{3/2} M_{pl} |c_\psi|^{5/2}. $$  \hspace{1cm} (3.41) 

[Recall $\Lambda_\omega$ is given by Eq. (3.27) and $M_* = M_B$.] One can take $c_1 \simeq M_{pl}$, but now a more reasonable choice is $c_1 \simeq M_*$. Then, the condition (3.41) becomes

$$ M_* < M_{pl} |c_\psi|^{1/2}, \quad (c_1 = M_*), $$  \hspace{1cm} (3.42) 

which is much less restricted than the one of $c_1 \simeq M_{pl}$. In addition, in order to have the sixth order derivative terms dominate, we must also require

$$ M_A \gtrsim M_* $$  \hspace{1cm} (3.43) 

Therefore, it is concluded that, provided that the conditions (3.41) and (3.43) hold, the HMT setup with any $\lambda$ [18, 21] does not have the strong coupling problem.

**IV. CONCLUSIONS**

In this paper, we have studied the strong coupling problem of a scalar field in the framework of the HMT setup [18], with an arbitrary coupling constant $\lambda$, generalized recently by da Silva [21]. As shown previously [20], when the energy of a process is higher than $\Lambda_\omega$, it becomes strong coupling. To avoid it, one can provoke the Blas-Pujolas-Sibiryakov (BPS) mechanism [20], in which a new energy scale $M_*$ is introduced, so that the sixth order derivative terms become important before the strong coupling energy scale $\Lambda_\gamma$ reaches. Once the high order derivative terms take over, the scaling behavior of the system is modified in such a way that all the nonrenormalizable terms become either strictly renormalizable or superrenormalizable, as shown explicitly in Sec. III.B.2. Whereby, the strong coupling problem is resolved.

It should be noted that, in order for the mechanism to work, $\lambda$ cannot be exactly one, as one can see from Eq. (3.41). In other words, the theory cannot reduce exactly to general relativity in the IR. However, since general relativity has achieved great success in low energies, $\lambda$ cannot be significantly different from one in the IR, in order for the theory to be consistent with observations. As first noticed by BPS in their model without projectability condition, the most stringent constraints come from the preferred frame effects due to Lorentz violation, which requires [20],

$$ |\lambda - 1| \lesssim 4 \times 10^{-7}, \quad M_* \lesssim 10^{15} \text{ GeV}. $$  \hspace{1cm} (4.1) 

In addition, the timing of active galactic nuclei [36] and gamma ray bursts [37] requires

$$ M_A \gtrsim 10^{10} \sim 10^{11} \text{ GeV}. $$  \hspace{1cm} (4.2) 

To obtain the constraint (4.1), BPS used the results from the Einstein-aether theory, as these two theories coincide in the IR [38].

In this paper, we have shown that the BPS mechanism is also applicable to the HMT setup. However, it is not clear whether the condition (4.1) is also applicable to the HMT setup, as the effects due to Lorentz violation in this setup have not been worked out, yet. On the other hand, the condition (4.2) is applicable, because this condition was obtained from the dispersion relations, which are the same in both setups.

In addition, the BPS mechanism cannot be applied to the Sotiriou-Visser-Weinfurter generalization with projectability condition [22], because the condition $M_* < \Lambda_\omega$, together with the one that instability cannot occur within the age of the universe, requires fine-tuning,

$$ |\lambda - 1| < 10^{-24}, $$  \hspace{1cm} (4.3) 

as shown explicitly in [27]. However, in the HMT setup with any $\lambda$, the Minkowski spacetime is stable [20], so such a fine-tuning does not exist.

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