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Constraints on scalar–tensor theory of gravity by the recent observational results on gravitational waves

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Abstract The speed of gravitational waves provides us a new tool to test alternative theories of gravity. The constraint on the speed of gravitational waves from GW170817 and GRB170817A is used to test some classes of Horndeski theory. In particular, we consider the coupling of a scalar field to Einstein tensor and the coupling of the Gauss–Bonnet term to a scalar field. The coupling strength of the Gauss–Bonnet coupling is constrained to be in the order of 10^{-15}. In the Horndeski theory we show that in order for this theory to satisfy the stringent constraint on the speed of GWs the mass scale M introduced in the non-minimally derivative coupling is constrained to be in the range 10^{15} GeV $\gg M \gtrsim 2\times10^{-35}$ GeV taking also under consideration the early times upper bound for the mass scale M. The large mass ranges require no fine-tuning because the effect of non-minimally derivative coupling is negligible at late times.

1 Introduction

The detection of gravitational waves (GWs) by the Laser Interferometer Gravitational-Wave Observatory (LIGO) Scientific Collaboration and Virgo Collaboration opens the window to study strong field gravitational physics and test alternative theories of gravity [1–6]. In particular, the recent detection of the GW170817 from the merger of a binary neutron star [6] and the electromagnetic counterparts starts a new era of multi-messenger GW astronomy. A gamma ray burst GRB170817A was observed 1.74 ± 0.05 s later by Fermi Gamma-Ray Burst Monitor [7] and the International Gamma-Ray Astrophysics Laboratory [8]. If we assume that the peak of the GW signal and the first photons were emitted simultaneously, and the 1.74 s time difference is caused by the faster speed of GWs, then we get an upper bound on the speed of GWs $c_{gw}/c - 1 \leq 7 \times 10^{-16}$ [9]. If we assume that the GRB signal was emitted 10s after the GW signal, then we get a lower bound $c_{gw}/c - 1 > -3 \times 10^{-15}$ [9]. The precise measurement of the propagation speed of GWs is a very powerful tool to test alternative theories of gravity [10–20].

Recently there is a lot of activity studying scalar–tensor theories [21] and one of them is the gravitational theory which is the result of the Horndeski Lagrangian [22]. Horndeski theories because they lead to second-order field equations can be technically simple, and they prove consistent without ghost instabilities [23]. In Horndeski theory the derivative self-couplings of the scalar field screen the deviations from GR at high gradient regions (small scales or high densities) through the Vainshtein mechanism [24], thus satisfying solar system and early universe constraints [25–32].

A subclass of Horndeski theories includes the coupling of the scalar field to Einstein tensor. This term introduces a new mass scale in the theory which on short distances allows to find black hole solutions [33–37], while a black hole can be formed if one considers the gravitational collapse of a scalar field coupled to the Einstein tensor [38]. On large distances the presence of the derivative coupling acts as a friction term in the inflationary period of the cosmological evolution [39–44]. Also, the preheating period at the end of inflation was studied, and it was found that there is a suppression of heavy particle production as the derivative coupling is increased. This was attributed to the fast decrease of kinetic energy of the scalar field because of its wild oscillations [45]. A holographic application was performed in [46] where it was shown that the change of the kinetic energy of the scalar field coupled to Einstein tensor allowed to holographically simulate the effects of a high concentration of impurities in a real material. The above discussion indicates that the coupling of the scalar field to Einstein tensor alters the kinematical properties of the scalar field.

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Assuming that the scalar field coupled to Einstein tensor plays the role of dark energy and drives the late cosmological expansion it was found [47,48] that the propagation speed of the tensor perturbations around the cosmological background with Friedmann–Robertson–Walker (FRW) metric is different from the speed of light c, so the measurement of the current speed of GWs can be used to test the applicability of this and Horndeski theories to explain the late time accelerated cosmological expansion [13–16,49–51]. Due to small deviation (on the order of 10^{-15}) of the current speed of GWs from the speed of light, in Refs. [13,14,16] it was argued that dark energy models that predict $c_{gw} \neq c$ at late cosmological times are ruled out, and in Horndeski theory the only viable nonminimal coupling to gravity has the conformal form $f(\phi)R$. However, with the help of derivative conformal or disformal transformations, some Horndeski and beyond Horndeski theories can survive the speed constraint [16,36,52,53].

In this work, we will perform a detailed analysis on the effect of the latest observational results on the current speed of GWs c_{gw} to the Horndeski theories with the non-minimally derivative coupling. From the early cosmological evolution we know that the derivative coupling of the scalar field to Einstein tensor alters the kinetic energy of the scalar field [40] influencing in this way the dynamical evolution of the Universe giving an upper bound to the mass scale coupling. For the late cosmological evolution we will perturb the FRW metric under tensor perturbations and we will show that a subclass of Horndeski theory consisting of the usual kinetic term and the coupling of the scalar field to Einstein tensor is still viable provided that the mass scale M introduced in the non-minimally derivative coupling is highly constrained from the recent results on the speed of GWs. The result also shows that no tuning on the parameter is needed to satisfy the stringent constraints on the speed of GWs. For comparison we also discuss bounds on the Gauss–Bonnet coupling from the observational bounds on c_{gw}.

The paper is organized as follows. In Sect. 2 we discuss the Horndeski theory, and studying the speed of tensor perturbations we obtain bounds on the mass scale introduced by the presence of the derivative coupling of the scalar field to Einstein tensor. In Sect. 3 we discuss the bounds on coupling α of the Gauss–Bonnet theory coupled to a scalar field. Finally, in Sect. 4 are our conclusions.

2 The effect of the speed of gravitational waves in the Horndeski theory

In this section we will briefly review the Horndeski theory, we will discuss the speed of GWs in this theory and assuming that the scalar field present in the Horndeski Lagrangian plays the role of dark energy we will find a lower bound on the mass scale introduced in derivative coupling of the scalar field to Einstein tensor.

The action of the Horndeski theory is given by [22],

$$S = \int d^4x\sqrt{-g}(L_2 + L_3 + L_4 + L_5), \tag{1}$$

where

$$\begin{aligned}
L_2 &= K(\phi, X), \quad L_3 = -G_3(\phi, X)\Box\phi,\\
L_4 &= G_4(\phi, X)R + G_{4,X}\left[(\Box\phi)^2 - (\nabla_\mu\nabla_\nu\phi)(\nabla^\mu\nabla^\nu\phi)\right],\\
L_5 &= G_5(\phi, X)G_{\mu\nu}\nabla^\mu\nabla^\nu\phi - \frac{1}{6}G_{5,X}[(\Box\phi)^3\\
&\quad -3(\Box\phi)(\nabla_\mu\nabla_\nu\phi)(\nabla^\mu\nabla^\nu\phi)\\
&\quad +2(\nabla^\mu\nabla_\alpha\phi)(\nabla^\alpha\nabla_\beta\phi)(\nabla^\beta\nabla_\mu\phi)],
\end{aligned}$$

with $X = -\nabla_\mu\phi\nabla^\mu\phi/2$, $\Box\phi = \nabla_\mu\nabla^\mu\phi$, the functions K, G_3, G_4 and G_5 are arbitrary functions of ϕ and X, and $G_{j,X}(\phi, X) = \partial G_j(\phi, X)/\partial X$ with $j = 4, 5$.

This action is the most general one for scalar–tensor theory with at most second-order field equations. If we take $K = G_3 = G_5 = 0$ and $G_4 = M_{\rm Pl}/2$, then we obtain Einstein's general relativity. If we take $G_3 = G_5 = 0$, $K = X - V(\phi)$, and $G_4 = f(\phi)$, then we get scalar–tensor $f(\phi)R$ theories. If we take $G_4 = M_{\rm Pl}^2/2 + X/(2M^2)$ or $G_4 = M_{\rm Pl}^2/2$ and $G_5 = -\phi/(2M^2)$, then we get the non-minimally derivative coupling $G_{\mu\nu}\nabla^\mu\phi\nabla^\nu\phi/(2M^2)$ with the mass scale M [54].

The stability of the Horndeski theory in the FRW background was studied in [55]. General conditions on the functions appearing in the Horndeski Lagrangian were given for the theory to be ghost free and stable under tensor perturbations. While in the flat background, the propagation speed of tensor perturbations is the same as the speed of light [56], in the cosmological FRW background the propagation speed of tensor perturbations in the Horndeski theory was found to be [54]

$$c_{gw}^2 = \frac{G_4 - X\left(\ddot{\phi}G_{5,X} + G_{5,\phi}\right)}{G_4 - 2XG_{4,X} - X\left(H\dot{\phi}G_{5,X} - G_{5,\phi}\right)}. \tag{2}$$

We are interested in the propagation speed of tensor perturbations of the subclass of the Horndeski theory that consists of the usual kinetic scalar field term and the coupling of the scalar field to Einstein tensor given by the action

$$S = \int d^4x\sqrt{-g}\left\{\frac{M_{\rm Pl}^2}{2}R - \frac{1}{2}\left[g_{\mu\nu} - \frac{1}{M^2}G_{\mu\nu}\right]\nabla^\mu\phi\nabla^\nu\phi - V(\phi)\right\}. \tag{3}$$

Perturbing the FRW metric as

$$ds^2 = -dt^2 + a^2(t)\left(\delta_{ij} + h_{ij}\right)dx^i dx^j, \tag{4}$$

and expanding the action (3) to the second order of the tensor perturbations h_{ij}, we obtain the quadratic action [44,48]

$$S = \frac{M_{\rm Pl}^2}{8}\int d^3x dt a^3\left[(1-\Gamma)\,\dot{h}_{ij}^2 - \frac{1}{a^2}(1+\Gamma)\left(\partial_k h_{ij}\right)^2\right], \tag{5}$$

where $\Gamma = \dot{\phi}^2/(2M^2M_{\rm Pl}^2)$. From the action (5) we derive the equation of motion for GWs

$$\ddot{h}_{ij} + 3H\dot{h}_{ij} - \frac{\dot{\Gamma}}{1-\Gamma}\dot{h}_{ij} - \frac{1+\Gamma}{1-\Gamma}\frac{\nabla^2 h_{ij}}{a^2} = 0. \tag{6}$$

Under the transverse-traceless gauge, the Fourier components of tensor perturbations $h_{ij}(\mathbf{x}, t)$ is

$$h_{ij}(\mathbf{x}, t) = \int d^3k\left[h_k^+(t)\,\varepsilon_{ij}^+ + h_k^\times(t)\,\varepsilon_{ij}^\times\right]\exp(i\mathbf{k}\cdot\mathbf{x}), \tag{7}$$

where $k^i\varepsilon_{ij}^s = \varepsilon_{ii}^s = 0$, $\varepsilon_{ij}^s\varepsilon_{ij}^{s'} = 2\delta_{ss'}$, and the superscript "$s$" stands for the "+" or "×" polarizations. Substituting Eqs. (7) into (6), we get

$$\ddot{h}_k^s + 3H\dot{h}_k^s\left[1 - \frac{\dot{\Gamma}}{3H(1-\Gamma)}\right] + \frac{c_{gw}^2k^2}{a^2}h_k^s = 0. \tag{8}$$

The propagation speed for both polarization states is

$$c_{gw}^2 = \frac{1+\Gamma}{1-\Gamma}. \tag{9}$$

This result can also be obtained from the general formula of the Horndeski theory given in Eq. (2) choosing $G_4 = M_{\rm Pl}^2/2 + X/(2M^2)$ or $G_4 = M_{\rm Pl}^2/2$ and $G_5 = -\phi/(2M^2)$. For the Horndeski theory, it was argued that the precise measurement on the speed of GWs $c_{gw} = 1$ requires $G_{4,X} = G_{5,\phi} = G_{5,X} = 0$, and only the conformal coupling $f(\phi)R$ is allowed [13,14,16]. At late times, since the effect of the non-minimally derivative coupling $G_{\mu\nu}\nabla^\mu\phi\nabla^\nu\phi \sim H^2\dot{\phi}^2/M^2$ is negligible compared with the canonical kinetic term $\dot{\phi}^2$ due to the decrease of the Hubble parameter H as the Universe expands, the speed given by Eq. (9) can be close to 1. Instead of requiring that $G_{4,X} = G_{5,\phi} = G_{5,X} = 0$, we show that a large mass range for the coupling M is allowed to satisfy the stringent constraint on the speed of GWs with negligible but nonzero deviation.

Using the the upper bound on the speed of of GWs [9]

$$\frac{c_{gw}}{c} - 1 \le 7\times 10^{-16}, \tag{10}$$

we obtain,

$$0 < \Gamma \le 7\times 10^{-16}. \tag{11}$$

This constraint is much less stringent than the classical constraint $\Gamma \le 2/3\times 10^{-20}$ derived from the constraint on the Parameterized Post-Newtonian (PPN) parameter $\alpha_3 = 6\Gamma < 4\times 10^{-20}$ [57].

Using the constraint (11) we will obtain a lower bound on the mass scale M of the derivative coupling. If we take the scalar field as dark energy and use the observational constraint $1 + w = \dot{\phi}^2/\rho_\phi = \dot{\phi}^2/(3M_{\rm Pl}^2H_0^2\Omega_\phi) \sim 0.2$ and $\Omega_m = 0.3$, then we get the contribution of the canonical kinetic energy as

$$\frac{(1+w)(1-\Omega_m)}{2} = \frac{\dot{\phi}^2}{6M_{\rm Pl}^2H_0^2} = \frac{\Gamma}{3}\frac{M^2}{H_0^2} \sim 0.07. \tag{12}$$

Combining Eqs. (11) and (12), we get the constraint on the coupling constant M

$$\frac{H_0^2}{M^2} \lesssim 3.3\times 10^{-15}. \tag{13}$$

For the theory with the coupling of a scalar field to Einstein tensor, the PPN parameters are [57]

$$\begin{aligned}&\beta = 1+6\Gamma,\ \ \gamma = 1+3\Gamma,\ \ \alpha_1 = 12\Gamma,\ \ \alpha_2 = 3\Gamma,\\ &\alpha_3 = 6\Gamma,\ \ \zeta_2 = 15\Gamma,\ \ \zeta_3 = 3\Gamma,\ \ \xi = \zeta_1 = \zeta_4 = 0.\end{aligned} \tag{14}$$

The stringent constraint coming from $\alpha_3 = 6\Gamma < 4\times 10^{-20}$ [58,59], substituting this result into Eq. (12), we get

$$\frac{H_0^2}{M^2} \lesssim 3.2\times 10^{-19}. \tag{15}$$

This constraint is much stronger than Eq. (13). If we use the constraint (13), we get the lower bound on the coupling M as $M \gtrsim 2\times 10^{-35}$ GeV. In the New Higgs inflation [47], the non-minimally derivative coupling enhances the friction of the expansion and the high friction limit requires $M \ll 10^{15}$ Gev [44] while in [60] limits on M are also discussed during the reheating period. Therefore, the mass scale introduced in the derivative coupling is 10^{15} GeV $\gg M \gtrsim 2\times 10^{-35}$ GeV.[1]

This result is interesting and it shows that the coupling of the scalar field to Einstein tensor has a complete different behaviour compared to a scalar field minimally coupled to gravity. While the kinetic energy of a minimally coupled scalar field practically does not understand the cosmological evolution, the kinetic energy of scalar field coupled to Einstein tensor changes as the Universe expands. At the inflationary epoch it can drive inflation with steep potentials while

[1] This lower bound of the mass M allows the sound speed squared of the tensor perturbations to reach the observational bounds of the GWs in the model discussed in [61] making in this way the model of unification of dark matter with dark energy in Horndeski theory viable.

as the Universe expands its contribution to the cosmological evolution is less important and at the late cosmological epoch is negligible, so GWs propagate at the speed of light at late times.

3 The effect of the speed of gravitational waves in the Gauss–Bonnet theory

Another non-trivial extension of GR which gives second order differential equations is the Lovelock theory [62], which apart from the Einstein–Hilbert term also includes higher order curvature terms. The simplest case is the Gauss–Bonnet theory which is a second order Lovelock theory which however in four dimensions is a topological invariance. If however it is coupled to a scalar field then a scalar–tensor theory is generated from the action [63],

$$S = \int \sqrt{-g} d^4x \left[\frac{M_{\rm Pl}^2}{2} R + X - V(\phi) + \frac{f(\phi)}{8} R_{GB} \right], \tag{16}$$

where

$$R_{\rm GB} = R^2 - 4R_{\mu\nu}R^{\mu\nu} + R_{\mu\nu\alpha\beta}R^{\mu\nu\alpha\beta}, \tag{17}$$

and we have also included a scalar potential $V(\phi)$.

It is interesting to notice that the Gauss–Bonnet theory coupled to a scalar field in four dimensions can be generated from the general Horndeski action (1) [54,64] making the following identifications of the functions involved

$$K(\phi, X) = X - V(\phi) + f''''(\phi) X^2 (3 - \ln X), \tag{18}$$

$$G_3(\phi, X) = \frac{f'''(\phi)}{2} X (7 - 3\ln X), \tag{19}$$

$$G_4(\phi, X) = \frac{M_{\rm Pl}^2}{2} + \frac{f''(\phi)}{2} X (2 - \ln X), \tag{20}$$

$$G_5(\phi, X) = -\frac{f'(\phi)}{2} \ln X. \tag{21}$$

Then we can use the general formula for the propagation of gravitational waves Eq. (2) for the functions (18)–(21) and we get

$$c_{gw}^2 = \frac{M_{\rm Pl}^2 + 2Xf''(\phi) + \ddot{\phi} f'(\phi)}{M_{\rm Pl}^2 + H\dot{\phi} f'(\phi)}. \tag{22}$$

The role that the Gauss–Bonnet tern coupled to a scalar field plays in the late cosmological evolution has been extensively studied [65–77]. In this work we will use a specific model discussed in [77]. In this model by choosing the scalar potential as

$$V(\phi) = \left[-\frac{\beta}{2} + \frac{1}{9}(3 - 8\alpha) \right] g(\phi) - \rho_0 e^{-\phi/\sqrt{\beta}}, \tag{23}$$

where α and β are model parameters. An exact solution of the gravitational field equations $\phi = 3\sqrt{\beta} \ln a$ was found with the coupling function of the scalar field to be

$$f(\phi) = \alpha \int \frac{8(3 - 4\alpha) d\phi}{3\sqrt{\beta} g(\phi)}, \tag{24}$$

where

$$g(\phi) = \frac{(3 - 4\alpha)A}{3} \exp\left(\frac{(8\alpha + 27\beta)\phi}{3\sqrt{\beta}(4\alpha - 3)} \right) + \frac{9(4\alpha - 3)\rho_0}{20\alpha + 27\beta - 9} e^{-\phi/\sqrt{\beta}}, \tag{25}$$

and A is an integration constant, ρ_0 is the present value of the energy density for matter. Note that if the coupling strength $\alpha = 0$ we do not have the Gauss–Bonnet coupling. From the solution, we get the current value of the ratio of the energy densities between dark energy and matter,

$$\frac{\Omega_{de}}{\Omega_m} = \frac{A(20\alpha + 27\beta - 9) - 3\rho_0(20\alpha + 27\beta)}{3\rho_0(20\alpha + 27\beta - 9)}, \tag{26}$$

and the current equation of state parameter for dark energy

$$w_{de} = -\frac{A(20\alpha + 27\beta - 9)^2}{3(4\alpha - 3)\left(c(20\alpha + 27\beta - 9) - 3\rho_0(20\alpha + 27\beta)\right)}. \tag{27}$$

If we take the current value of the ratio Ω_{de}/Ω_m to be 7/3 [78], then we obtain the integration constant A,

$$A = \frac{(200\alpha + 270\beta - 63)\rho_0}{20\alpha + 27\beta - 9}, \tag{28}$$

and the current equation of state parameter for dark energy

$$w_{de} = -\frac{63 - 200\alpha - 270\beta}{63 - 84\alpha}. \tag{29}$$

Using Eq. (22) the speed of GWs is

$$c_{gw}^2 = 1 - \frac{320\alpha^2 + 6\alpha(90\beta - 11)}{45 + 160\alpha^2 - 180\alpha}. \tag{30}$$

Using the bound on the speed of GWs [9]

$$-3 \times 10^{-15} \leq \frac{c_{gw}}{c} - 1 \leq 7 \times 10^{-16}, \tag{31}$$

and the constraint $-1.1 < w_{de} < -0.9$ [78], we get the range on the model parameters α and β as shown in Fig. 1.

The results in Fig. 1 show that for a range of values of the parameter β, the coupling strength of the Gauss–Bonnet

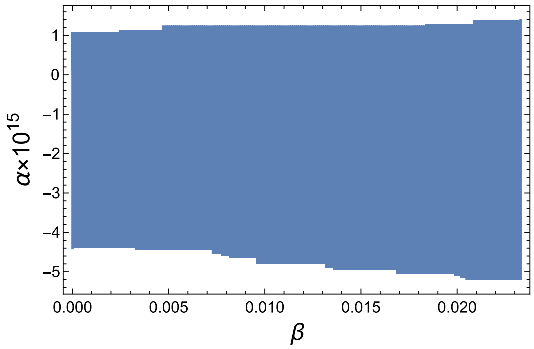

Fig. 1 The constraint on the coupling constant α and the model parameter β

term is in the order of 10^{-15}. For the inflationary model with $V(\phi) \sim \phi^2/2$ and $f(\phi) = -8\alpha\phi^2/2$, the absolute value of the coupling strength α is constrained to be less than the order of 0.01 [72]. For the power-law inflation with both exponential coupling and potential, the coupling strength is constrained to be $-1\times10^{-4} < \alpha < 4\times10^{-4}$ [73]. Therefore, the constraint from the speed of GWs is much stronger.

4 Conclusions

The first measurement of the speed of GWs by GW170817 and GRB170817A bounds the deviation of the speed of GWs from the speed of light to be no more than one part in 10^{15}, so it provides the evidence that $c_{gw} = c$. Using these observational result we can test alternative theories of gravity for their validity to describe the cosmological evolution at late times.

We used these bounds on c_{gw} to constrain first a subclass of the Horndeski theory in which a scalar field except its minimal coupling is also coupled to Einstein tensor. Assuming that the scalar field plays the role of dark energy we found a lower bound on the mass scale introduced by this coupling and combining the constraints from inflation the energy scale of the derivative coupling is bounded to be 10^{15} GeV $\gg M \gtrsim 2\times10^{-35}$ GeV. This result requires no fine-tuning and shows that it is possible to get $c_{gw} \approx c$ from the terms with $G_{4,X} \neq 0$ and $G_{5,\phi} \neq 0$ if their effects are negligible at late times.

We also studied the Gauss–Bonnet theory in four dimensions coupled to a scalar field. The coupling of the Gauss–Bonnet term to scalar field not only gives successful inflation, but also provides late time cosmic acceleration. Using a particular model with a specific form of the coupling function $f(\phi)$ which allows an exact solution of the gravitational equations, we found that the bounds on c_{gw} constrains the Gauss–Bonnet coupling strength to be $\alpha \lesssim 10^{-15}$, a constraint much stronger than the coupling strength $-1\times10^{-4} < \alpha < 4\times10^{-4}$ resulted from models with power-law inflation.

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