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► To cite this version:

Alexandru Dima, Filippo Vernizzi. Vainshtein Screening in Scalar-Tensor Theories before and after GW170817: Constraints on Theories beyond Horndeski. 2018. cea-01697808

HAL Id: cea-01697808

<https://cea.hal.science/cea-01697808>

Preprint submitted on 31 Jan 2018

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Vainshtein Screening in Scalar-Tensor Theories before and after GW170817: Constraints on Theories beyond Horndeski

Alexandru Dima¹ and Filippo Vernizzi²

¹*SISSA – International School for Advanced Studies, Via Bonomea 265, 34136 Trieste, Italy*

²*Institut de physique théorique, Université Paris Saclay CEA, CNRS, 91191 Gif-sur-Yvette, France*

(Dated: December 14, 2017)

Screening mechanisms are essential features of dark energy models mediating a fifth force on large scales. We study the regime of strong scalar field nonlinearities, known as Vainshtein screening, in the most general scalar-tensor theories propagating a single scalar degree of freedom. We first develop an effective approach to parameterize cosmological perturbations beyond linear order for these theories. In the quasi-static limit, the full nonlinear effective Lagrangian contains six independent terms, one of which starts at cubic order in perturbations. We compute the two gravitational potentials around a spherical body. Outside and near the body, screening reproduces standard gravity, with a modified gravitational coupling. Inside the body, the two potentials are different and depend on the density profile, signalling the breaking of the Vainshtein screening. We provide the most general expressions for these modifications, revising and extending previous results. We apply our findings to show that the combination of the GW170817 event, the Hulse-Taylor pulsar and stellar structure physics, constrain the parameters of these general theories at the level of a few $\times 10^{-1}$, and of GLPV theories at the level of 10^{-2} .

Introduction.— The recent simultaneous observation of gravitational waves and gamma ray bursts from GW170817 [1] and GRB 170817A [2] has allowed to constrain very precisely the relative speed between gravitons and photons. This measurement has had dramatic impact on the parameter space of modified gravity theories characterized by a single scalar degree of freedom [3–6]. In particular, the so-called Horndeski theories [7, 8], a class of well-studied scalar-tensor theories that are often used as benchmarks to parameterize modifications of gravity, have been drastically simplified. Their higher-order Lagrangian terms, quadratic and cubic in second derivatives of the field, predict a speed of gravitational waves that differ from that of light and are thus ruled out. This fact has triggered renewed interest for the surviving theories, i.e. those extending the Horndeski class that are compatible with the GW170817 observation, such as certain subclasses of GLPV theories [9, 10].

Lagrangian terms with higher-derivatives are crucial to suppress, via the so-called Vainshtein mechanism [11, 12], the fifth force exchanged by the scalar and responsible for the modifications of gravity on large scale. On the other hand, theories extending the Horndeski class are known to display a breaking of the Vainshtein screening inside matter [13], a phenomenon that has allowed to constrain the parameter space of these theories with astrophysical observations [14–18].

The purpose of this paper is to study the Vainshtein mechanism in the general framework of the degenerate theories introduced in [19, 20], which includes the Horndeski class and theories beyond Horndeski [9, 10, 21]. We will consider only theories that can be related to the Horndeski class by an invertible metric redefinition [20, 22]. In the classification of Ref. [19], they are called degenerate higher-order scalar-tensor theories of Type Ia. Moreover, we will focus on theories up to *quadratic* in the second derivative of the scalar field. In particular, we do not consider the *cubic* theories [23], whose Vainshtein mechanism has been poorly studied due to its complexity, because they are anyway ruled out by the observation of GW170817 [3, 5, 6].

We do so by reducing these theories to their essential elements with the use of the Effective Field Theory of dark energy description developed in [24–29]. Moreover, we focus on scales much smaller than the Hubble radius and we restrict to non-relativistic sources, in which case the scalar fluctuations satisfy the quasi-static approximation. We will first derive very general expressions for the two potentials in the Vainshtein regime. This will be important to extend and clarify previously obtained results. Then, restricting to theories propagating gravitons at the speed of light, we will use our expressions to show that a combination of constraints from stellar structure [15, 30] and from precise measurements of the decrease of the orbital period in the Hulse-Taylor binary pulsar [31] severely constrain these scenarios.

During the preparation of this work, Refs. [32, 33] have appeared, where some of the results derived in this article are independently obtained using different approaches.

Degenerate Higher-Order Scalar-Tensor theories.— Let us consider a tensor-scalar field theory described by an action including all possible quadratic combinations up to second derivatives of the field ϕ [19],

$$S = \int d^4x \mathcal{L} = \int d^4x \sqrt{-g} \left[P(\phi, X) + Q(\phi, X) \square \phi + f(\phi, X) {}^{(4)}R + \sum_{I=1}^5 a_I(\phi, X) L_I(\phi, \phi_{;\nu}, \phi_{;\rho\sigma}) \right], \quad (1)$$

where ${}^{(4)}R$ is the 4D Ricci scalar. A semicolon denotes the covariant derivation, $X \equiv -\phi_{;\mu} \phi^{;\mu}/2$, and the L_I are defined by

$$L_1 = \phi_{;\mu\nu} \phi^{;\mu\nu}, \quad L_2 = (\phi_{;\mu}^{;\mu})^2, \quad L_3 = (\phi_{;\mu}^{;\mu})(\phi^{;\rho} \phi_{;\rho\sigma} \phi^{;\sigma}), \\ L_4 = \phi^{;\mu} \phi_{;\mu\nu} \phi^{;\nu\rho} \phi_{;\rho}, \quad L_5 = (\phi^{;\rho} \phi_{;\rho\sigma} \phi^{;\sigma})^2. \quad (2)$$

In the following, we are going to focus on Type Ia theories, which satisfy $a_1 + a_2 = 0$ and two other degeneracy conditions

[19] that fix two functions, for instance a_4 and a_5 . Degenerate theories that are not in this class have been shown to propagate scalar fluctuations with sound speed squared with opposite sign to the sound speed squared of tensor fluctuations [28] and we will not consider them here. The theory with $a_3 + a_4 = a_5 = 0$ is degenerate also in the absence of gravity. In particular, this includes the case $f = G_4$, $a_1 = -a_2 = -G_{4,X}$ (a comma denotes the derivative with respect to the argument) and $a_3 = a_4 = a_5 = 0$, corresponding to quartic Horndeski theories, and the case $f = G_4$, $a_1 = -a_2 = -G_{4,X} + 2XF_4$, $a_3 = -a_4 = -2F_4$ and $a_5 = 0$, corresponding to quartic GLPV theories. The functions P and Q do not affect the degeneracy character of the theory.

Effective Theory of Dark Energy.— To describe cosmological perturbations around a FRW solution in theories with a preferred slicing induced by a time-dependent scalar field, it is convenient to use the EFT of dark energy. To formulate the action (1) with the conditions $a_1 + a_2 = 0$, we use the ADM metric decomposition, where the line element reads $ds^2 = -N^2 dt^2 + h_{ij}(dx^i + N^i dt)(dx^j + N^j dt)$, and we choose the time as to coincide with the uniform field hypersurfaces. Moreover, we are going to focus only on the operators that contribute to the *quasi-static* limit.

In this gauge, expanded around a FRW background, the full nonlinear action reads

$$S_{\text{QS}} = \int d^4x \sqrt{h} \frac{M^2}{2} (-\delta \mathcal{K}_2 + c_T^2 {}^{(3)}R + 4H\alpha_B \delta K \delta N + (1 + \alpha_H) {}^{(3)}R \delta N + \alpha_V \delta N \delta \mathcal{K}_2 + 4\beta_1 \delta K V + \beta_2 V^2 + \beta_3 a_i a^i). \quad (3)$$

Here $H \equiv \dot{a}/a$ (a dot denotes the time derivative), $\delta N \equiv N - 1$, $\delta K_i^j \equiv K_i^j - H \delta_i^j$ is the perturbation of the extrinsic curvature of the time hypersurfaces, δK its trace and ${}^{(3)}R$ is the 3D Ricci scalar of these hypersurfaces. Moreover, $\delta \mathcal{K}_2 \equiv \delta K^2 - \delta K_i^j \delta K_j^i$, $V \equiv (\dot{N} - N^i \partial_i N)/N$ and $a_i \equiv \partial_i N/N$.

We have also defined the effective Planck mass, which normalizes the graviton kinetic energy, by $M^2 \equiv 2(f - 2a_2 X)$ and a few independent parameters, related to the functions in (1) by

$$\begin{aligned} \alpha_B &= \alpha_V + \dot{\phi}(f_\phi + 2Xf_{,\phi X} + XQ_{,X})/(M^2 H), \\ c_T^2 &= 2f/M^2, \quad \alpha_H = 4X(a_2 - f_{,X})/M^2, \\ \beta_1 &= 2X(f_{,X} - a_2 + a_3 X)/M^2, \\ \alpha_V &= 4X(f_{,X} - 2a_2 - 2Xa_{2,X})/M^2. \end{aligned} \quad (4)$$

The function c_T^2 is the fractional difference between the speed of gravitons and photons. Sometimes called braiding [34], the function α_B measures the kinetic mixing between metric and scalar fluctuations [24]. The function α_H measures the kinetic mixing between matter and the scalar in GLPV theories [9, 10, 35] and vanishes for Horndeski theories. The functions $\beta_{1,2,3}$ parameterize the presence of higher-order operators. In the EFT of dark energy formulation, the degeneracy conditions that ensure that the action (3) describes the propagation of a single scalar degree of freedom are [28]

$$\beta_2 = -6\beta_1^2, \quad \beta_3 = -2\beta_1 [2(1 + \alpha_H) + \beta_1 c_T^2], \quad (5)$$

	M^2	α_B	$c_T^2 - 1$	α_V	α_H	β_1
$P(\phi, X)$	0	0	0	0	0	0
$Q(\phi, X) \square \phi$	0	✓	0	0	0	0
Quartic Horndeski	✓	✓	✓	✓	0	0
Quartic GLPV	✓	✓	✓	✓	✓	0
Quadratic DHOST	✓	✓	✓	✓	✓	✓
After GW170817	free	free	0	$-\alpha_H$	free	free

TABLE I. Lagrangian operators of the EFT of dark energy allowed in various theories and the consequences of the equality between speed of gravity and light on these theories.

so that we do not need the explicit expression for β_2 and β_3 in terms of the functions defining (1). We will impose these conditions later. Finally, the operator proportional to α_V is the only one that starts cubic in the perturbations. In the nonlinear EFT action, it was introduced (as $-\alpha_{V1}$) in [29] to describe nonlinear dark energy perturbations. Notice that the action (3) does not include the kinematicity [34] Lagrangian term $\alpha_K \delta N^2$, because it can be neglected in the quasi-static limit [29]. The total number of independent parameters, and thus of Lagrangian operators, is thus six.

We summarize the relation between the EFT operators and the corresponding covariant Lagrangians in Table I, where we also state in which way these operators are affected by the equality between the speed of gravity and light, see [3] and discussion below.

Action in Newtonian gauge.— We now expand the Lagrangian (1) around an FRW background. We consider only scalar fluctuations in the Newtonian gauge, where $\delta N = \Phi$ and $h_{ij} = a(t)^2(1 - 2\Psi)\delta_{ij}$ and $N^i = 0$. Without loss of generality, we take $\phi = t + \pi(t, \vec{x})$.

In the quasi-static regime, time derivatives are of order Hubble and the Lagrangian is dominated by terms with $2(n - 1)$ spatial derivatives for n fields. Considering only these terms, one obtains

$$\begin{aligned} S_{\text{QS}} &= \int d^4x \frac{M^2 a}{2} \left[(c_1 \Phi + c_2 \Psi + c_3 \pi) \nabla^2 \pi + c_4 \Psi \nabla^2 \Phi \right. \\ &\quad + c_5 \Psi \nabla^2 \Psi + c_6 \Phi \nabla^2 \Phi + (c_7 \dot{\Psi} + c_8 \dot{\Phi} + c_9 \dot{\pi}) \nabla^2 \pi \\ &\quad + \frac{b_1}{a^2} \mathcal{L}_3^{\text{Gal}} + \frac{1}{a^2} (b_2 \Phi + b_3 \Psi) \mathcal{E}_3^{\text{Gal}} + \frac{1}{a^2} (b_4 \nabla_i \Psi + b_5 \nabla_i \Phi \\ &\quad \left. + b_6 \nabla_i \dot{\pi}) \nabla_j \pi \Pi^{ij} + \frac{1}{a^4} (d_1 \mathcal{L}_4^{\text{Gal}} + d_2 \nabla_i \pi \nabla_j \pi [\Pi^2]^{ij}) \right]. \end{aligned} \quad (6)$$

Here, adopting the notation of [13], $\Pi_{ij} \equiv \nabla_i \nabla_j \pi$, $\Pi_{ij}^n \equiv \nabla_i \nabla^{k_1} \pi \nabla_{k_1} \nabla^{k_2} \pi \dots \nabla^{k_{n-1}} \nabla_j \pi$ and $[\Pi^n] \equiv \delta^{ij} \Pi_{ij}^n$ we have defined $\mathcal{L}_3^{\text{Gal}} \equiv -\frac{1}{2}(\nabla \pi)^2 [\Pi]$, $\mathcal{L}_4^{\text{Gal}} \equiv -\frac{1}{2}(\nabla \pi)^2 \mathcal{E}_3^{\text{Gal}}$ and $\mathcal{E}_3^{\text{Gal}} \equiv [\Pi]^2 - [\Pi^2]$. The coefficients c_i , b_i and d_i are time-dependent functions related to the functions P , Q , f , a_i defining (1), and their derivatives, evaluated on the background solution.

Equivalently, eq. (6) can also be obtained from the EFT action (3), after introducing the scalar fluctuation π by a time

diffeomorphism $t \rightarrow t + \pi(t, \vec{x})$ [36]. In this case, the coefficients c_i , b_i and d_i can be expressed in terms of the EFT parameters. The coefficients $c_{1,2,3}$ and b_1 are functions of M^2 , α_B , c_T^2 and H (and their time derivatives) but we do not need their explicit expressions for the following discussion. The other coefficients are given by

$$\begin{aligned} c_4 &= 4(1 + \alpha_H), \quad c_5 = -2c_T^2, \quad c_6 = -\beta_3, \\ c_7 &= 4\alpha_H, \quad c_8 = -2(2\beta_1 + \beta_3), \quad c_9 = 4\beta_1 + \beta_3, \\ b_2 &= \alpha_V - \alpha_H - 4\beta_1, \quad b_3 = c_T^2 - 1, \quad b_4 = -c_7, \\ b_5 &= -c_8, \quad b_6 = -2c_9, \quad d_1 = -b_3 - b_2, \quad d_2 = c_9. \end{aligned} \quad (7)$$

The relevant nonlinear couplings dominating in the Vainshtein regime will be the quartic ones in (6), i.e. those proportional to d_1 and d_2 . Note that they contain c_T^2 , α_H , α_V , β_1 and β_3 , but not α_B . See more on this below. For $\beta_{1,2,3} = 0$, it is straightforward to verify that the above action agrees with those given in [37] for Horndeski and in [13] for GLPV theories.

To study the behaviour of Φ , Ψ and π around dense matter sources, we add to the action (6) the coupling with non-relativistic matter with energy density $\rho_m = \bar{\rho}_m(t) + \delta\rho_m(t, \vec{x})$, i.e.,

$$S_m = - \int d^4x a^3 \Phi \delta\rho_m. \quad (8)$$

Vainshtein mechanism.— To study the Vainshtein regime, we take matter to be described by some overdensity, spherically distributed around the origin. We define

$$x \equiv \frac{1}{\Lambda^3} \frac{\pi'}{a^2 r}, \quad y \equiv \frac{1}{\Lambda^3} \frac{\Phi'}{a^2 r}, \quad z \equiv \frac{1}{\Lambda^3} \frac{\Psi'}{a^2 r}, \quad \mathcal{A} \equiv \frac{1}{8\pi M \Lambda^3} \frac{m}{r^3}, \quad (9)$$

where a prime denotes the derivative with respect to the radial distance r , $m(t, r) \equiv 4\pi \int_0^r \tilde{r}^2 \delta\rho_m(t, \tilde{r}) d\tilde{r}$ and Λ is some mass scale of order $\Lambda \sim (MH^2)^{1/3}$. Integrating over space the equations obtained by varying the action (6) respectively with respect to Φ and Ψ , and using Stokes theorem, we obtain

$$(c_1 - \dot{c}_8 - 3Hc_8)x + 2c_6y + c_4z - c_8\dot{x} + 2\Lambda^3x[(2b_2 - b_5)x - b_5rx'] = 2M\mathcal{A}, \quad (10)$$

$$(c_2 - \dot{c}_7 - Hc_7)x + c_4y + 2c_5z - c_7\dot{x} + 2\Lambda^3x[(2b_3 - b_4)x - b_4rx'] = 0. \quad (11)$$

By applying the analogous procedure to the equation obtained by varying the action with respect to π , we get

$$\begin{aligned} &2\tilde{c}_3x + \tilde{c}_1y + \tilde{c}_2z + 2\tilde{c}_9\dot{x} + c_8\dot{y} + c_7\dot{z} + 2c_9\ddot{x} \\ &+ 2\Lambda^3\{2\tilde{b}_1x^2 + (5Hb_6 + \dot{b}_6)rxx' + b_6(5x\dot{x} + 2rx\dot{x}' + r\dot{x}x') \\ &+ x[(4b_2 + 3b_5)y + (4b_3 + 3b_4)z + b_5ry' + b_4rz']\} \\ &+ 8\Lambda^6\{(d_1 + 3d_2)x^3 + d_2x[r^2(x')^2 + rx(6x' + rx'')]\} = 0. \end{aligned} \quad (12)$$

The coefficients with the tildes are related to those without the tildes and to their time derivatives, but we do not need their explicit expression for what follows. Notice that these equations contains terms with up to second derivatives (in time and

space), indicating that the equations of motion are higher than second order.

Equations (10) and (11) are linear in y and z , and can be solved for these two variables and their solutions can be replaced in eq. (12) to obtain an equation for x only. Using the definitions (7) for the time-dependent coefficients c_i , b_i and d_i and imposing the degeneracy conditions (5), the space and time derivatives on x cancel and one remains with

$$x^3 + v_1x^2 + (v_2 + v_3\mathcal{A} + v_4\mathcal{A}')x + v_5\mathcal{A} + v_6\mathcal{A}' = 0, \quad (13)$$

where the coefficients v_i are related to the original EFT functions (4). The fact that x in this equation always appears without derivatives is not surprising because the theory is degenerate and the scalar degree of freedom must satisfy second-order equations of motion. For large overdensities $\mathcal{A}$, this equation can be solved by

$$x^2 \approx -v_3\mathcal{A} - v_4\mathcal{A}' \quad (\mathcal{A} \gg 1), \quad (14)$$

which can be used to solve eqs. (10) and (11) for y and z in the regime where π is nonlinear. Imposing the degeneracy condition (5) on these coefficients, one finally obtains

$$\begin{aligned} \Phi'(r) &= G_N \left(\frac{m(r)}{r^2} + \gamma_1 m''(r) \right), \\ \Psi'(r) &= G_N \left(\frac{m(r)}{r^2} + \gamma_2 \frac{m'(r)}{r} + \gamma_3 m''(r) \right). \end{aligned} \quad (15)$$

These expressions give the two gravitational potentials close to the matter source, in the regime of large scalar field nonlinearities. Outside the matter source $m' = m'' = 0$ and one recovers the Newtonian behaviour, $\Phi = \Psi = G_N m/r$, although the coupling constant G_N is in general time-dependent and affected by α_V and β_1 . To parameterize this possible deviation from standard gravity in screened regions, we introduce the function

$$\gamma_0 \equiv (8\pi G_N M^2)^{-1} = 1 + \alpha_V - 3\beta_1. \quad (16)$$

Inside the matter source, the gravitational potentials are in general different and receive corrections that depend on the density profile of the source (and its radial derivative), similarly to what happens in beyond Horndeski theories [13]. The corrections are proportional to three time-dependent functions parameterizing the breaking of the Vainshtein screening inside matter, which can be expressed in terms of the parameters c_T^2 , α_H , α_V and β_1 as

$$\begin{aligned} \gamma_1 &\equiv \frac{(\alpha_H + c_T^2 \beta_1)^2}{c_T^2(1 + \alpha_V - 4\beta_1) - \alpha_H - 1}, \\ \gamma_2 &\equiv -\frac{\alpha_H(\alpha_H - \alpha_V + 2(1 + c_T^2)\beta_1) + \beta_1(c_T^2 - 1)(1 + c_T^2\beta_1)}{c_T^2(1 + \alpha_V - 4\beta_1) - \alpha_H - 1}, \\ \gamma_3 &\equiv -\frac{\beta_1(\alpha_H + c_T^2 \beta_1)}{c_T^2(1 + \alpha_V - 4\beta_1) - \alpha_H - 1}. \end{aligned} \quad (17)$$

The above expressions, derived here for the first time, are the most general for scalar-tensor theories propagating a single

scalar degree of freedom. We now focus on specific cases and compare with results previously found in the literature.

Beyond Horndeski theories.— Let us specialize eq. (17) to the beyond-Horndeski (or GLPV) theories, which do not contain higher derivatives in the EFT action (3), i.e. $\beta_{1,2,3} = 0$. In this case, the expressions for the γ_i simplify, i.e. $\gamma_0 = 1 + \alpha_V$,

$$\gamma_1 = \frac{\alpha_H^2}{c_T^2(1 + \alpha_V) - \alpha_H - 1}, \quad \gamma_2 = -\frac{\alpha_H(\alpha_H - \alpha_V)}{c_T^2(1 + \alpha_V) - \alpha_H - 1}, \quad (18)$$

and $\gamma_3 = 0$.

These equations extend the expressions obtained in [14, 15] under the assumption that $Q = f_\phi = 0$, in the notation of eq. (1). The expressions in those references are analogous to the ones above but with α_B replacing α_V . At first, it is surprising that α_B appears in those expressions because, in contrast with α_V , in the quasi-static limit the operator proportional to α_B does not contain terms quartic in the perturbations—such as the last two terms in eq. (6)—and that hence contribute to the Vainshtein mechanism. The explanation is that when $Q = f_\phi = 0$ one sees from eq. (4) that $\alpha_V = \alpha_B$, so that these expressions can also be written in terms of α_B . However, in general γ_1 and γ_2 are independent of α_B and one needs to go beyond the quadratic action and introduce the dependence on α_V .

After GW170817.— The simultaneous observation of GW170817 and GRB 170817A implies that gravitational waves travel at the speed of light, with very small deviations [38]. This has dramatically constrained the parameter space of available theories. To avoid that even 10^{-5} fluctuations in the matter overdensity and gravitational potentials along the path of the gravitons affect their speed of propagation, some of the coefficients of the EFT action must be simply set to zero.

It is straightforward to derive these consequences using eq. (3). The coefficient $c_T^2 - 1$ detunes the space kinetic term of the gravitons, $(\partial_k \gamma_{ij})^2$, contained in R , from their time kinetic term, $\dot{\gamma}_{ij}^2$, contained in $\delta\mathcal{H}_2$. Thus, one must require $c_T = 1$. This must be true also for small fluctuations around the background solution. Small fluctuations of the background induce small changes in δN in front of $\delta\mathcal{H}_2$ and R , which modify the speed of gravitons if the two coefficients α_H and $-\alpha_V$ do not coincide. Therefore, gravitons travel at the same speed as photons independently of small changes in the background if [3]

$$c_T = 1, \quad \alpha_V = -\alpha_H. \quad (19)$$

In eq. (4), this translates into $a_2 = a_{2,X} = 0$. Remarkably, these conditions are stable under quantum corrections [3, 39].

Using these conditions in eq. (17) we find $\gamma_0 = 1 - \alpha_H - 3\beta_1$ and

$$\gamma_1 = -\frac{(\alpha_H + \beta_1)^2}{2(\alpha_H + 2\beta_1)}, \quad \gamma_2 = \alpha_H, \quad \gamma_3 = -\frac{\beta_1(\alpha_H + \beta_1)}{2(\alpha_H + 2\beta_1)}. \quad (20)$$

These expressions, also obtained in [32, 33], are in general valid independently of α_B . For $\beta_1 = 0$ and $Q = f_\phi = 0$, they

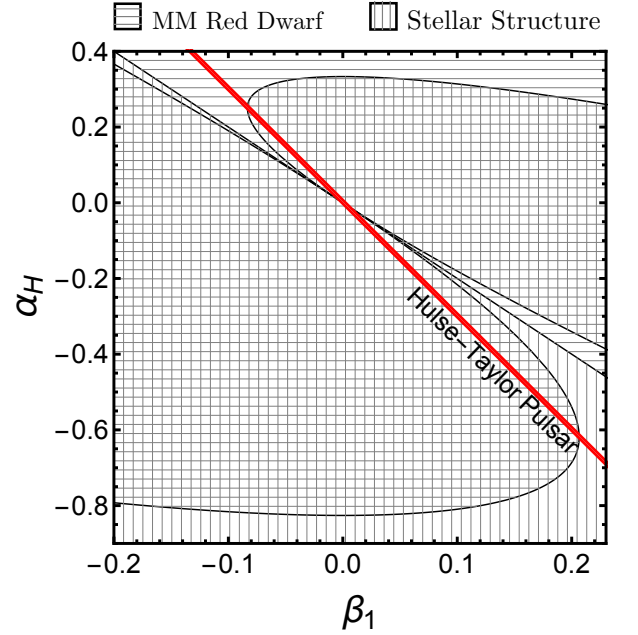


FIG. 1. Allowed region in the plane (β_1, α_H) after the GW170817 event, from the stellar structure condition $\gamma_1 > -1/6$ [15] (vertical grid), the minimal mass (MM) red dwarf $\gamma_1 < 0.4$ (v3 of [30]) (horizontal grid—the region with squares is allowed by both) and the Hulse-Taylor pulsar $-7.5 \times 10^{-3} \leq \gamma_0 - 1 \leq 2.5 \times 10^{-3}$ at 2σ (red band) [31]. The region allowed by the three observations is given by the overlap of the red band and the region with squares.

agree with [4] where $\alpha_B = -\alpha_H$ and become $\gamma_1 = -\alpha_H/2$, $\gamma_2 = \alpha_H$ and $\gamma_3 = 0$.

Before discussing the observational constraints on these expressions, we note that the second condition in eq. (19) does not necessarily apply if dark energy has a fixed ϕ independent of H , in which case small changes around the background do not induce a change in δN [3], and if (dark) matter is not coupled to the same metric as photons [29]. In these cases one must use the general expressions in eq. (17).

Observational constraints.— Several late-time observational bounds have been put on the parameters γ_i . The Newtonian potential Φ controls the stellar structure equation, so one can bound γ_1 independently of γ_2 and γ_3 . A negative value of γ_1 means stronger gravity inside a star so that, for stars to exist in hydrostatic equilibrium, one requires $\gamma_1 > -1/6$ [15]. An upper bound, $\gamma_1 < 0.4$, comes from requiring that the smallest observed red dwarf star has a mass larger than the minimum mass allowing hydrogen to burn in stars, see v3 of Ref. [30]. Constraints on γ_2 and γ_3 require observations involving the curvature potential Ψ and we do not discuss them here.

Let us turn now to γ_0 , i.e. the ratio between the screened effective Newton constant, G_N , and the effective coupling constant for gravitons, M^{-2} . As shown in Ref. [31], the decrease of the orbital period of binary stars is proportional to $(M^2 G_N c_T)^{-1}$. With $c_T = 1$, γ_0 can be constrained by the 40 year-long observation of the Hulse-Taylor pulsar (PSR B1913+16) [40]. Using the results of this reference, one ob-

tains $-7.5 \times 10^{-3} \leq \gamma_0 - 1 \leq 2.5 \times 10^{-3}$ at 2σ . This constraint assumes that the scalar radiation does not participate to the energy loss. For cubic screening, the effect has been shown to be suppressed by $-3/2$ powers of the product of the orbital period and the Vainshtein radius [41, 42].

As shown in Fig. 1, combining these constraints places tight bounds on α_H and β_1 : only the tiny overlap between the red band and the region with squares survives, which shows that $-0.59 \leq \alpha_H \leq 0.26$ and $-0.08 \leq \beta_1 \leq 0.20$. For GLPV theories ($\beta_1 = 0$), this leads to a very stringent bound on α_H : $-2.5 \times 10^{-3} \leq \alpha_H \leq 7.5 \times 10^{-3}$ at 2σ .

Conclusions.— We have obtained the most general expression of the gravitational potentials in the Vainshtein regime, for degenerate higher-order scalar-tensor theories up to quadratic in second derivatives of the scalar. After GW170817, these are the most general Lorentz-invariant theories propagating a single scalar degree of freedom. To do so, we have employed the EFT of dark energy approach at nonlinear order and computed the deviations from general relativity, outside and inside matter sources. In general, these modifica-

tions imply four observational parameters and depend on four EFT parameters, but if gravitons travel at the same speed as photons independently of small changes in the background, they depend only on α_H and β_1 , which measure the beyond-Horndeski “character” of the theories, see Table I. Using these results and those from astrophysical observations, we have obtained stringent constraints on these two parameters. Our bounds have been derived using only γ_0 and γ_1 . Any constraints on γ_2 and γ_3 will exclude a new region of the (β_1, α_H) plane, possibly ruling out theories beyond Horndeski.

Acknowledgements.— We thank D. Langlois, K. Noui and F. Piazza and especially P. Creminelli for enlightening discussions and suggestions and K. Koyama, J. Sakstein and M. Vallisneri for useful correspondences. Moreover, we thank N. Bartolo, P. Karmakar, S. Matarrese and M. Scomparin for kindly sharing the draft of their paper [43] before submission. A.D. acknowledges kind hospitality at IPHT during the completion of this work. F.V. acknowledges financial support from “Programme National de Cosmologie and Galaxies” (PNCG) of CNRS/INSU, France and the French Agence Nationale de la Recherche under Grant ANR-12-BS05-0002.

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