Speed of gravitational waves and the fate of scalar-tensor gravity

Dario Bettoni,[1,*] Jose María Ezquiaga,[2,†] Kurt Hinterbichler,[3,‡] and Miguel Zumalacárregui[1,4,§]

[1]Nordita, KTH Royal Institute of Technology and Stockholm University, Roslagstullsbacken 23, SE-106 91 Stockholm, Sweden

[2]Instituto de Física Teórica UAM/CSIC, Universidad Autónoma de Madrid, C/ Nicolás Cabrera 13-15, Cantoblanco, Madrid 28049, Spain

[3]CERCA, Department of Physics, Case Western Reserve University, 10900 Euclid Avenue, Cleveland, Ohio 44106, USA

[4]Berkeley Center for Cosmological Physics, LBNL and University of California at Berkeley, Berkeley, California 94720, USA

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The direct detection of gravitational waves (GWs) is an invaluable new tool to probe gravity and the nature of cosmic acceleration. A large class of scalar-tensor theories predicts that GWs propagate with velocity different than the speed of light, a difference that can be $\mathcal{O}(1)$ for many models of dark energy. We determine the conditions behind the anomalous GW speed, namely, that the scalar field spontaneously breaks Lorentz invariance and couples to the metric perturbations via the Weyl tensor. If these conditions are realized in nature, the delay between GW and electromagnetic signals from distant events will run beyond human time scales, making it impossible to measure the speed of GWs using neutron star mergers or other violent events. We present a robust strategy to exclude or confirm an anomalous speed of GWs using eclipsing binary systems, the electromagnetic phase of which can be exquisitely determined. The white dwarf binary J0651 + 2844 is a known example of such a system that can be used to probe deviations in the GW speed as small as $c_g/c - 1 \gtrsim 2 \times 10^{-12}$ when LISA comes online. This test will either eliminate many contender models for cosmic acceleration or wreck a fundamental pillar of general relativity.

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I. INTRODUCTION AND SUMMARY

The direct detection of gravitational radiation [1,2] has initiated a new era for astronomy, astrophysics, and fundamental physics. The observed gravitational wave (GW) events and the ones to come will usher in novel ways to test the nature of gravity [3]. Here, we will argue that probing the speed of GWs will be a decisive test for gravity and dark energy models.

The nature of the propagation of GWs is a question of great and fundamental interest. Einstein's General Relativity (GR) predicts two massless tensor polarizations, each traveling at the speed of light, c, with an amplitude inversely proportional to the distance from the source [4]. However, major outstanding theoretical issues such as the nature of dark energy and dark matter have led researchers to consider the possibility that gravity differs from GR in some regimes (see, e.g., Refs. [5,6] for reviews). In alternative theories of gravity, additional polarizations may propagate, each with potentially different velocities, attenuations, and effective masses [7]. This issue has been well studied in cosmology and has been a topic of discussion in connection to the early [8–11] and the late Universe [12–15]. There are fairly model-independent tests for effects caused by additional polarizations [16], damping [17–19], mass [20], and Lorentz symmetry violations [21,22]. To date, the speed of GWs has been upper bounded with the arrival timing of GW150914 between the two LIGO detectors [23]. Also, it has been constrained at the ∼1% level from the variation of the orbital period in binary pulsars [24]. Moreover, if $c_g < c$, a very stringent lower bound $c_g/c - 1 \gtrsim -10^{-15}$ can be obtained from the absence of gravitational Cherenkov radiation, as probed by ultrahigh-energy cosmic rays [25,26].

In this paper, we analyze the speed of GWs, c_g, in generic scalar-tensor theories of gravity and ask when it can differ from the speed of light, c. Unlike previous studies, we do not assume a specific cosmological background, instead focusing on the local speed of gravity. Such anomalous propagation is potentially observable if both gravitational waves and an electromagnetic (EM) or other nongravitational counterpart signal can be seen from the same source.

One of two scenarios will arise. The simultaneous arrival of a GW signal with a nongravitational counterpart from a distant source will set extremely stringent and model-independent bounds on c_g. However, a very slight difference in propagation speed (as predicted by many models of cosmic acceleration) would cause a delay between the signals' arrivals much larger than the multimessenger observation campaign. In this case, a GW signal never gets identified with its true EM counterpart, and other techniques

*dario.bettoni@nordita.org
†jose.ezquiaga@uam.es
‡kurt.hinterbichler@case.edu
§miguelzuma@berkeley.edu

must be used. We will discuss one such method, the phase lag test with eclipsing binaries, based on monitoring periodic galactic sources observable in GWs by future space missions such as LISA [27], and in EM by other means, and comparing the phase of the two signals.

A measurement of nontrivial c_g would have profound implications for our understanding of gravity. As we shall see, the anomalous propagation of GWs is directly related to fundamental properties of the underlying gravitation theories, which can hence be distinguished on this basis. Conversely, an observation consistent with GWs traveling at the speed of light will place much more severe constraints than any other available test on the large class of theories predicting an anomalous GW speed. In fact, current cosmological constraints on general scalar-tensor theories are only of the order of $\mathcal{O}(1-0.5)$ [28], while future forecasts will reach $\mathcal{O}(0.1-0.01)$ [29]. Testing the speed of GWs will dramatically improve these constraints to $\mathcal{O}(10^{-12}\text{–}10^{-17})$.

II. SCALAR FIELDS AND THE SPEED OF GWs

In the following, we are going to present a general method to compute the speed of GWs. Let us start with an example theory that predicts anomalous GWs propagation: a quartic shift-symmetric Horndeski theory [30,31] $S=\int d^4x\sqrt{-g}\mathcal{L}$ with

$$\mathcal{L}=G(X)R+G'(X)((\Box\phi)^2-\nabla_\mu\nabla_\nu\phi\nabla^\mu\nabla^\nu\phi), \quad (1)$$

where $X\equiv-\frac{1}{2}(\partial\phi)^2$ and $G'\equiv\partial G/\partial X$. We set $c=1$ in this section. Expanding around a background solution, $g_{\mu\nu}\to g_{\mu\nu}+h_{\mu\nu}$, $\phi\to\phi+\varphi$, yields a quadratic action for the fluctuations

$$\mathcal{L}=\frac{1}{2}[h_{\mu\nu}\mathcal{D}^{\mu\nu,\rho\sigma}h_{\rho\sigma}+h_{\mu\nu}\mathcal{D}^{\mu\nu}\varphi+\varphi\mathcal{D}\varphi], \quad (2)$$

where $\mathcal{D}^{(\cdots)}$ represent differential operators depending on the background fields $g_{\mu\nu}$ and ϕ and their derivatives.

Since we are interested in local propagation, we adopt Riemann normal coordinates around a point P and expand the scalar and metric background in a Taylor series about P, $g_{\mu\nu}=\eta_{\mu\nu}-\frac{1}{3}R_{\mu\rho\nu\sigma}x^\rho x^\sigma+\cdots$, $\phi=\phi_0+\phi_\mu x^\mu+\frac{1}{2}\phi_{\mu\nu}x^\mu x^\nu+\cdots$, where $\phi_\mu=\nabla_\mu\phi$, $\phi_{\mu\nu}=\nabla_\mu\nabla_\nu\phi$ and the derivatives and curvatures are all evaluated at P. This leaves freedom for a rotation and boost around P.

We may now zoom in and obtain an effective action valid around the point P by taking the scaling limit, $\lambda\to0$, with

$$x^\mu\to\lambda x^\mu, \qquad \varphi\to\frac{1}{\lambda}\varphi, \qquad h_{\mu\nu}\to\frac{1}{\lambda}h_{\mu\nu}. \quad (3)$$

The result is a flat space action, depending on the background field values and derivatives evaluated at P.

We will focus on the spin-2 polarizations present in GR and neglect the additional scalar mode. Imposing the transverse gauge condition $\partial^\mu h_{\mu\nu}=0$, the scaling-limit action reads

$$\mathcal{L}=\frac{1}{2}h_{\mu\nu}[G\Box+G'\phi^\rho\phi^\sigma\partial_\rho\partial_\sigma]h^{\mu\nu}+h^\rho_\mu G'\phi^\mu\phi^\nu\Box h_{\nu\rho}+\cdots, \quad (4)$$

where we omitted terms involving both the trace of the metric and the scalar field. We then perform a standard $3+1$ split of $h_{\mu\nu}$ and restrict to the transverse-traceless (TT) part of the spatial metric components h_{ij},

$$h_{00}=0, \qquad h_{0i}=0, \qquad h_{ij}=h^{TT}_{ij}, \qquad \varphi=0, \quad (5)$$

with $\partial^j h^{TT}_{ij}=\delta^{ij}h^{TT}_{ij}=0$. We will further assume that the spatial shear of the background scalar configuration is negligible.[1] This assumption simplifies the analysis, ensuring that h^{TT}_{ij} decouple from the other perturbations and allowing us to ignore the terms omitted in Eq. (4), which describe the scalar polarization and nondynamical metric elements.

If the field gradient ϕ_μ is timelike (as expected for a cosmological contribution), we can rotate the coordinates so that $\phi_\mu=(\dot\phi,0,0,0)$, for some constant $\dot\phi$. Then, the last term of (4) does not contribute, and

$$\mathcal{L}=\frac{1}{2}\{[G-G'\dot\phi^2](\dot h^{TT}_{ij})^2-G(\vec\nabla h^{TT}_{ij})^2\}, \quad (6)$$

from which we can read off the propagation speed

$$c_g^2=\frac{1}{1-\frac{G'}{G}\dot\phi^2}. \quad (7)$$

In particular, GR corresponds to $G(X)=$ const., and we recover $c_g=1$.

In the case of a spacelike field gradient, we can boost our reference frame so that the time component vanishes. Decomposing the gradient in components parallel and perpendicular to the GW propagation, $\phi_i=\phi_i^\parallel+\phi_i^\perp$, we obtain that the velocity of propagation of GWs depends on the direction as

$$c_g^2=1+\frac{G'|\phi_\parallel|^2}{G+G'|\phi_\perp|^2}. \quad (8)$$

In general, the speed is anisotropic (i.e., direction dependent), and equal for both the $+$ and $\times$ GW polarizations.

The scaling limit (3) eliminates all the lower derivative terms, which is the reason why the resulting GW speed is frequency independent. This is different for other well-studied cases, such as massive gravitons [32] (see Refs. [33,34] for reviews) or Lorentz violations. These other scenarios modify the waveform in a frequency-dependent way and can thus be constrained from GW

[1]The precise condition is $\phi_{ii}-\phi_{jj},\phi_{ij}\ll G'/G$ for $(i\neq j)$. This is satisfied in a boosted frame with $\phi_i=0$ whenever ϕ_μ is timelike.

observations alone [3,35,36]. For the sake of simplicity, we have also neglected the scalar mode, which may also have its own anomalous propagation speed [37–40].

III. CONDITIONS FOR ANOMALOUS GWs SPEED

We now study the origin of the anomalous speed of GWs (7), (8) in more generality. The Lagrangian for the transverse-traceless components (6) can be written in terms of an *effective gravitational metric*,

$$\mathcal{L} \propto h^{TT}_{\alpha\beta}(\mathcal{G}_{\mu\nu}\partial^\mu\partial^\nu)h^{\alpha\beta}_{TT}, \tag{9}$$

determining the causal structure of GW propagation.[2] The propagation path for GWs will be given by the condition $\mathcal{G}_{\mu\nu}dx^\mu dx^\nu = 0$ and will in general be different from the light-cone condition $g_{\mu\nu}dx^\mu dx^\nu = 0$ unless the two metrics obey a *conformal relation*: $\mathcal{G}_{\mu\nu} = \Omega(x)g_{\mu\nu}$. The lack of proportionality is found already in the simple example theory (1), where

$$\mathcal{G}_{\mu\nu} = G(X)g_{\mu\nu} + G'(X)\phi_\mu\phi_\nu \tag{10}$$

and $\mathcal{G}_{\mu\nu}$ and $g_{\mu\nu}$ are connected by a *disformal relation* [41] for which $\mathcal{G}_{\mu\nu} \neq \Omega(x)g_{\mu\nu}$. Such a relation is ubiquitous in modern scalar-tensor theories [42–45].

Let us examine the conditions for a disformal relation to arise in a generic theory of gravity. First, it is necessary that the background scalar field has a nontrivial configuration that spontaneously breaks Lorentz invariance, e.g., $\phi_\mu \neq 0$ in Eq. (10). In addition, we note that the effective second-order Lagrangian (2) follows from the second variation of the action over a background and is hence equal to the first variation of the equations of motion (EoM). The simplest term in the EoM producing second derivatives and entering in Eq. (9) is the Ricci curvature. When expanded to first order, considering only the TT components,

$$R^{TT}_{\mu\nu} = -\frac{1}{2}\Box h^{TT}_{\mu\nu} \quad \text{and} \quad R^{TT} = 0 \tag{11}$$

only contribute to the conformal part in the effective gravitational metric (9).

Further, second derivative terms are restricted by covariance to originate either from the Riemann tensor or repeated application of covariant derivatives (e.g., third derivatives of the scalar field), with the two cases related by $\nabla_\mu\nabla_\nu\phi^\alpha = \nabla_\nu\nabla_\mu\phi^\alpha + R^\alpha_{\lambda\mu\nu}\phi^\lambda$. To first order, the TT contribution to the Riemann tensor reads

$$R^{TT}_{\mu\alpha\nu\beta} = -\frac{1}{2}\partial_\beta\partial_\alpha h^{TT}_{\mu\nu} + \frac{1}{2}\partial_\nu\partial_\alpha h^{TT}_{\mu\beta} - (\alpha \leftrightarrow \mu). \tag{12}$$

The above expression explicitly induces disformal terms in Eq. (9) via contractions with scalar field derivatives. In the simple example (1), only ϕ^μ enters in the effective metric (10) due to the particular nonminimal coupling to the Ricci scalar. In more general cases, for instance, when there are couplings to the Ricci tensor such as in quintic Horndeski, second derivatives $\phi^{\mu\nu}$ could appear contracted with the derivatives of the metric and hence in $\mathcal{G}_{\mu\nu}$. Thus, the effective metric would belong to the extended disformal class [43,46]. In any case, because the Ricci tensor only contributes to the conformal part, the contribution of $R_{\mu\nu\alpha\beta}$ leading to the anomalous speed of GWs is fully captured by the Weyl tensor (i.e., the trace-free part of the Riemann tensor). For the simple theory (1), the Weyl tensor appears explicitly in the equations of motion whenever $G' \neq 0$ [47].

These considerations allow us to formulate a *Weyl criterion* for anomalous speed of spin-2 GWs. The effective gravitational metric of the example theory (10) can be generalized to

$$\mathcal{L} \propto h_{\mu\nu}(\mathcal{C}\Box + \mathcal{W}^{(\alpha\beta)}\partial_\alpha\partial_\beta)h^{\mu\nu}, \tag{13}$$

where $\mathcal{C}$ and $\mathcal{W}^{\mu\nu}$ are the contributions associated with the Ricci and Weyl tensors, respectively. Anomalous GW speed requires that $\mathcal{W}^{\alpha\beta} \neq 0$, i.e., for the background scalar derivatives to couple to the Riemann/Weyl curvature. If the Weyl factor is purely timelike and constant around P, $\mathcal{W}^{\mu\nu} = \mathcal{W}^{00}\delta^\mu_0\delta^\nu_0$, the speed of tensors becomes

$$c_g^2 = \frac{\mathcal{C}}{\mathcal{C} - \mathcal{W}^{00}}. \tag{14}$$

In Horndeski theories, which are a general framework that encompasses most of the current dark energy models, the EoM are second order [30]. Therefore, the occurrence of the Weyl tensor fully distinguishes theories in which $c_g = c$ exactly and those in which the speed of GWs is allowed to vary. GR, kinetic gravity braiding [48], and Jordan-Brans-Dicke theories [49] (including $f(R)$ [50,51]) only contain Ricci curvature in their equations of motion and therefore do not modify the speed of GWs. On the other hand, covariant Galileons [52] and the covariantization of other generalizations [53–56] will generically predict $c_g \neq c$ [57].

Although the Weyl criterion is characteristic of Scalar-tensor (ST) theories, the occurrence of a disformal relation can be applied to more general theories such as massive gravity [32]. In this case, the kinetic term has the Einstein-Hilbert form and hence $c_g = c$ plus corrections $\mathcal{O}(\frac{m^2}{E^2})$ beyond the scaling limit (3), as expected from unbroken Lorentz invariance. In the case of bigravity [58], the situation is more subtle, as the kinetic term of the second metric $\sqrt{-f}R[f_{\mu\nu}]$ forces its excitations to propagate along $f_{\mu\nu}dx^\mu dx^\nu = 0$, with $f_{\mu\nu} \neq \Omega(x)g_{\mu\nu}$ in nonflat background space-times. Although matter

[2] We focus on the spin-2 components and assume they decouple. Nonetheless, Eq. (9) remains valid for the propagation eigenstates of the linearized fields (including the scalar mode and the generalization of $h^{TT}_{\alpha\beta}$ when it couples to other perturbations), with a different $\mathcal{G}^A_{\mu\nu}$ for each polarization A.

does not couple to $f_{\mu\nu}$ directly, the anomalous speed may be detectable via graviton oscillations [59,60], as well as in doubly coupled theories [61]. Many theories that attempt to explain away dark matter such as TeVeS also predict an anomalous GW speed [62].

IV. PHASE LAG TEST WITH ECLIPSING BINARIES

Most of the present bounds on c_g can be significantly strengthened by comparing GWs with other signals. In theories in which matter is universally coupled to the metric, electromagnetic signals and ultrarelativistic particles propagate at the speed of light. This produces a delay between GW and electromagnetic signals,

$$\Delta t = r\left(\frac{1}{c_g} - \frac{1}{c}\right) \equiv \frac{r}{c}\epsilon_g \approx 10^{14} s \frac{r}{\mathrm{Mpc}} \epsilon_g, \tag{15}$$

where we define the *differential delay parameter* $\epsilon_g \equiv c\partial\Delta t/\partial r$ (in general space-times, r is the proper distance, and one has to correct for time dilation at emission [16]). The detection of violent, multimessenger events at cosmological distances bears the promise of phenomenal constraints, even in the presence of considerable astrophysical uncertainties. LIGO expects to perform such measurements using violent events such as binary compact object mergers involving neutron stars [63].

However, no distant GW-EM event will possibly be observed if c_g is modified significantly, since the delay between both signals will be much larger than the monitoring time around the GW detection. This is the case of cosmic acceleration models without a cosmological constant such as covariant Galileons [52,64], for which $|c_g/c - 1| \sim 10$–100% (see Ref. [57] and Fig. 1 of Ref. [65]). If such a model is responsible for cosmic acceleration, the arrival times of both signals will differ by millions or even billions of years. Clearly, an alternative test for the speed of GWs would be needed in this situation. In the following, we discuss how observations of sources with periodic signals can help to test whether $c_g = c$. In particular, we propose a *phase lag test with eclipsing binaries* that overcomes this limitation.

The anomalous speed of GWs can be tested by monitoring periodic sources with both GW and EM emission [66,67]. This ensures that both signals can be observed continuously and allows for a long observation period. A suitable source is a binary system in the band of space-based interferometers [68], including *verification binaries* [69–71]: systems expected to be resolvable by LISA and that have already been identified and characterized using electromagnetic observations (see Ref. [71] for an updated list). An extraordinarily clean binary system is WDS J0651 + 2844: a binary, detached white dwarf system ~1 kpc away from the Sun and of which orbital plane is

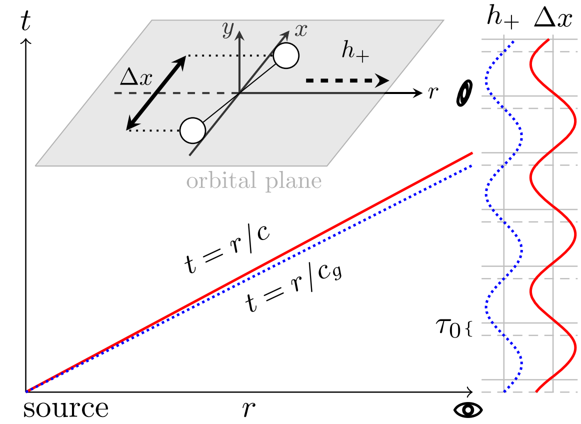

FIG. 1. The phase lag test for the speed of gravity. A compact binary system such as WDS J0651 + 2844 is monitored both electromagnetically and using GWs. For this geometry (top), only the + GW polarization is emitted in the observer's direction. Its amplitude h_+ is initially correlated with the object transverse separation Δx, but a phase lag (16) accumulates on the propagation if $c_g \neq c$ (bottom and right).

approximately aligned with the Solar System, allowing the observation of periodic eclipses [72]. Its short orbital period ~12.75 min falls within the LISA band and makes it a loud GW source, in which the effect of GW emission has already been observed by the period variation [73].

Let us model WDS J0651 + 2844 as a binary orbit coplanar with the observer and at a distance r from it, cf. Fig. 1. Because of symmetry, the gravitational radiation emitted in the observer's direction will be predominantly in the + polarization $h_{ij} = h_+(t)(\hat{x}\hat{x} - \hat{y}\hat{y})$.[3] Assuming GR (i.e., $c_g = c$), the h_+ polarization will be in phase with Δx, the distance between the objects transverse to the line of sight as observed electromagnetically. Therefore, although the components of the binary will not be resolvable, $\Delta x = 0$ coincides with the eclipses and can be timed with extraordinary precision [66].

In theories other than GR, the EM and GW observables will evolve as periodic functions of different retarded times, i.e., $\Delta x \propto \cos(2\omega(t - r/c))$ and $h_+ \propto \cos(2\omega(t - r/c_g))$. The difference in propagation speed accumulated over the propagation distance r produces a *phase lag* between the GW and the EM signals,[4]

$$\Delta\Phi(t) = 2\omega\frac{r(t)}{c}\left(\frac{c}{c_g} - 1\right) = 2\omega\frac{r(t)}{c}\epsilon_g, \tag{16}$$

where the distance between the source and detector

[3] The orbital inclination is $\iota = 86.9^{+1.6}_{-1.0}$ deg [72], making $h_\times$ suppressed by $\cos(\iota) \approx 0.05$ in amplitude and shifted $\pi/2$ in phase relative to the + component.

[4] We have neglected the delay from the atmospheric or interstellar refractive index, which can be shown to be unimportant [66].

$$r(t) = r_0 + v_{\rm rel}t + r_{\rm orb}(t) \tag{17}$$

includes the initial separation, relative velocity, and detector's orbit. We will focus on the effect of $r_0, v_{\rm rel}$, as the effect of $r_{\rm orb}$ has been considered [74].

For eclipsing binaries, we can neglect the error in EM measurements in constructing the relative phase (16) $\Delta\Phi(t) \equiv 2\omega(\tau_0 + \hat{\beta}t)$. The precision will be then limited by our knowledge of the GW signal. We can obtain an estimate of the $1-\sigma$ uncertainties using the Fisher matrix formalism [75] for the quantities

$$\tau_0 \equiv \varepsilon_g \frac{r_0}{c}, \quad \Delta\tau_0 = \frac{1}{\sqrt{2}\omega\Sigma} \approx 0.2s\left(\frac{2\pi/\omega}{765s}\right)\left(\frac{T}{5y}\right), \tag{18}$$

$$\hat{\beta} \equiv \varepsilon_g \frac{v_{\rm rel}}{c}, \quad \Delta\hat{\beta} = \frac{\sqrt{3/2}}{\omega T\Sigma} \approx 10^{-8}\left(\frac{2\pi/\omega}{765s}\right)\left(\frac{T}{5y}\right), \tag{19}$$

where T is the observation time and Σ denotes the total signal-to-noise ratio of the GW detection (see the Appendix). The expected detection significance of verification binaries with LISA is $\Sigma \sim 100(\frac{T}{1y})$ [74].

A nonzero measurement of either Eq. (18) or Eq. (19) represents a smoking gun for $c_g \neq c$:

(i) τ_0: The relative phase of the signals can detect an anomalous propagation speed in the range $|\varepsilon_g| \gtrsim 2\times 10^{-12}(\frac{\rm kpc}{r_0})(\frac{\Delta\tau_0}{0.2s})$. The false-negative case in which $2r_0\varepsilon_g\omega/(c\pi)$ equals an integer within the measurement error is very unlikely (probability $\approx \Sigma^{-1} \sim 0.2\%$) and can be excluded by observing multiple systems or measuring the frequency shift $\hat{\beta}$.

(ii) $\hat{\beta}$: The relative velocity of the system induces a frequency shift, sensitive to anomalous GW speeds in the range $|\varepsilon_g| \gtrsim 10^{-4}(\frac{30\ {\rm km/s}}{v_{\rm rel}})(\frac{\Delta\beta}{10^{-8}})$. Despite the $(\omega T)^{-1}$ gain when observing over many cycles, this test is less competitive due to the nonrelativistic factor.

Note that both the measurement of the relative phase and the velocity can be used as a test of $\varepsilon_g \neq 0$ *and* as a measurement of c_g. The latter application requires a measurement of either r_0 or $v_{\rm rel}$, which will almost certainly dominate the error. Nevertheless, clean systems such as WDS J0651 + 2844 will be able to confirm deviations from $c_g = c$ at the level of few parts in a trillion.

V. CONCLUSIONS

Many well-studied models of dark energy and modified gravity theories predict an anomalous local speed of gravity around nontrivial backgrounds. The Weyl criterion provides a clear-cut way to distinguish two classes of gravitational theories, those for which the speed of GWs is *exactly equal* to the speed of light and those in which it can vary depending on the theory parameters and the background configuration of the scalar field. Future multimessenger GW observations will probe this effect to exquisite precision: if the prediction of GR is satisfied, this will place such a stringent constraint on theories allowing variations in the speed of GWs, $\mathcal{O}(10^{-17})$, that they will become uninteresting for any low-energy application, including cosmic acceleration. On the other hand, a confirmation of an anomalous propagation of GWs by extragalactic and galactic sources would be able to rule out GR and all other theories with simple kinetic terms, which would significantly impact our understanding of gravity. This could be achieved applying the proposed phase lag test for eclipsing binaries to the already identified white dwarf binary WDS J0651 + 2844. Either of these two scenarios shows that the speed of GWs will be by far one of the most powerful tools to constrain gravity and dark energy models.

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APPENDIX: SIGNAL-TO-NOISE ESTIMATES

The signal-to-noise ratio Σ for a GW detection is given by

$$\Sigma^2 = \frac{1}{\sigma_f^2}\int_0^T \tilde{R}^2(t)dt \equiv \varrho. \tag{A1}$$

Here, $\tilde{R}$ is the response of the detector to the signal, and σ_f^2 is the noise power at the GW frequency. We assume the GW to be monochromatic and follow Ref. [74] (see Ref. [75] for further details and cautionary notes). For a given detector, the response function depends on the GW polarizations as $\tilde{R}(t) = A_+(t)h_+ + A_\times(t)h_\times$, where A_i contain information about the antenna pattern of the detector and its orientation as a function of time. However, as discussed in the text, we will consider the situation in which only one polarization is received and assume that the errors in the electromagnetic signal are negligible. Therefore, we can reconstruct the relative phase [Eq. (16) in the main text] directly,

$$\tilde{R}(t) = \Upsilon\cos(\varpi t + \psi), \tag{A2}$$

where the signal has an overall amplitude Υ, which will not directly affect the reconstruction of ψ and ω.

The Fisher matrix is then given as the derivative of Eq. (A1) with respect to the model parameters

$$F_{ij} = \frac{2}{\sigma_f^2}\int_0^T \frac{\partial\tilde{R}}{\partial\theta_i}\frac{\partial\tilde{R}}{\partial\theta_j}dt, \tag{A3}$$

where $\theta_i = (\Upsilon, \varpi, \psi)$ collectively denotes the unknown parameters of the signal. The error in the parameter θ_i assuming the other ones are perfectly known is $(F_{ii})^{-1/2}$, while the error in a parameter marginalized over the rest is $\sqrt{(F^{-1})_{ii}}$.

The Fisher matrix elements read

$$
\begin{aligned}
F_{\Upsilon\Upsilon} &= \frac{2}{\sigma_f^2}\int \cos^2(\varpi t + \psi)dt = 2\varrho/\Upsilon^2,\\
F_{\Upsilon\varpi} &= \frac{2}{\sigma_f^2}\int -t\sin(\varpi t + \psi)\Upsilon\cos(\varpi t + \psi)dt \sim \text{osc.},\\
F_{\Upsilon\psi} &= \frac{2}{\sigma_f^2}\int -\Upsilon\cos(\varpi t + \psi)\sin(\varpi t + \psi)dt \sim \text{osc.},\\
F_{\varpi\varpi} &= \frac{2}{\sigma_f^2}\int \Upsilon^2 t^2 \sin^2(\varpi t + \psi)dt = 2\varrho\frac{t^2}{3} + \text{osc.},\\
F_{\varpi\psi} &= \frac{2}{\sigma_f^2}\int \Upsilon^2 t \sin^2(\varpi t + \psi)dt = \varrho t + \text{osc.},\\
F_{\psi\psi} &= \frac{2}{\sigma_f^2}\int \Upsilon^2 \sin^2(\varpi t + \psi)dt = 2\varrho + \text{osc.},
\end{aligned}
$$

where osc. denotes oscillatory terms that become negligible for $T \gg \varpi^{-1}$ and we have used $\varrho = \frac{\Upsilon^2}{2\sigma_f^2}T$. Since $F_{\Upsilon\varpi}, F_{\Upsilon\psi}$ do not build up with time, the amplitude is uncorrelated with the frequency and the phase. However, ϖ and ψ are correlated with one another. The Fisher matrix and its inverse for the (ϖ, ψ) subspace are

$$
\begin{aligned}
\hat{F} &= \varrho\begin{pmatrix} \frac{2}{3}T^2 & T \\ T & 2 \end{pmatrix},\\
\hat{F}^{-1} &= \frac{1}{\varrho}\begin{pmatrix} \frac{6}{T^2} & -\frac{3}{T} \\ -\frac{3}{T} & 2 \end{pmatrix},
\end{aligned} \tag{A4}
$$

from which we read the errors in the phase and frequency,

$$
\begin{aligned}
\Delta\psi &= \frac{\sqrt{2}}{\Sigma},\\
\Delta\varpi &= \frac{\sqrt{6}}{T\cdot\Sigma},
\end{aligned} \tag{A5}
$$

which translate straightforwardly into the results [Eqs. (18) and (19) in the main text].

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