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Spontaneous scalarization of black holes and compact stars from a Gauss–Bonnet coupling

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We identify a class of scalar-tensor theories with coupling between the scalar and the Gauss–Bonnet invariant that exhibit spontaneous scalarization for both black holes and compact stars. In particular, these theories formally admit all of the stationary solutions of general relativity, but these are not dynamically preferred if certain conditions are satisfied. Remarkably, black holes exhibit scalarization if their mass lies within one of many narrow bands. We find evidence that scalarization can occur in neutron stars as well.

Introduction. Gravitational wave observations [1–7] allow us to probe the structure of black holes (BHs) with unprecedented accuracy. Hence, they can reveal the existence of new fundamental scalar fields [8, 9], provided that they leave an imprint on BHs. However, no-hair theorems (see [10, 11] for reviews) dictate that conventional scalar-tensor theories will have the same stationary, asymptotically flat BH solutions as general relativity (GR) [12–14]. In spherical symmetry [15] and slow rotation [16, 17], this result extends to generalized scalar-tensor theories, i.e. theories that exhibit derivative self-interactions and derivative couplings between the scalar and curvature invariants, provided that the scalar respects shift symmetry.

One could still detect scalars in these theories through the imprint they leave when they are excited [18, 19]. One can also circumvent no-hair theorems by violating some of their assumptions [20–23]. No-hair theorems also help single out particularly interesting theories that have hairy BHs. A well-studied example is the action

$$S = \frac{1}{2}\int \mathrm{d}^4x\sqrt{-g}\left[R - \frac{1}{2}\nabla_\alpha\varphi\nabla^\alpha\varphi + f(\varphi)\mathscr{G}\right] + S_\mathrm{m}[g_{\mu\nu},\psi]\,, \tag{1}$$

where $\mathscr{G} \equiv R_{\mu\nu\rho\sigma}R^{\mu\nu\rho\sigma} - 4R_{\mu\nu}R^{\mu\nu} + R^2$ is the Gauss-Bonnet invariant. We use geometrical units with $c = 8\pi G = 1$ and the mostly plus metric signature. The scalar field φ is coupled to $\mathscr{G}$, which has dimensions of length^{-4} ($\equiv L^{-4}$), through a function $f(\varphi)$, with dimensions L^2. The matter fields ψ are minimally coupled to the metric $g_{\mu\nu}$ through the action S_m. We will refer to this class of theories as scalar-Gauss-Bonnet (sGB) gravity. When f is exponential the theory is well-known to admit hairy BHs [24], whereas a linear f yields the only shift-symmetric theory with second-order field equations that exhibits BH hair [16, 17] (despite the no-hair theorem of [15]).

The main purpose of this paper is to demonstrate that a new subclass of theories, contained in (1), exhibits a particularly interesting phenomenon: BH spontaneous scalarization. As we demonstrate below, this subclass of theories generically admits solutions where the scalar field is constant and the metric satisfies Einstein's equations. However, under certain conditions these solutions are unstable, and solutions where the scalar field in nontrivial are dynamically preferred. This leads to hairy BHs only when the BH mass lies within certain ranges. Compact stars in these theories also exhibit spontaneous scalarization. The mechanism resembles that proposed by Damour and Esposito-Farèse [25], where there is a coupling between φ and the trace of the stress-energy tensor, T. However, there are important differences – most notably the fact that the effect is present for BHs as well.

A no-hair theorem in sGB and how to evade it. We start by identifying the class of theories in question. Varying (1) with respect to φ and $g_{\mu\nu}$ yields

$$\Box\varphi = -f_{,\varphi}\mathscr{G}\,, \tag{2a}$$

$$R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R = T_{\mu\nu}\,. \tag{2b}$$

Here $T_{\mu\nu}$ is the sum of the matter stress-energy tensor $T^\mathrm{m}_{\mu\nu} \equiv -(2/\sqrt{-g})(\delta S_\mathrm{m}/\delta g^{\mu\nu})$, plus a contribution coming from the variation of the φ-dependent part of the action with respect to the metric (see e.g. [24]).

Eq. (2a) does not admit φ = constant solutions, unless

$$f_{,\varphi}(\varphi_0) = 0\,, \tag{3}$$

for some constant φ_0. We consider Eq. (3) as an *existence condition* for GR solutions and focus on theories that satisfy it. This excludes the widely studied class of dilatonic theories where $f \sim \exp(\varphi)$ and the shift-symmetric $f \sim \varphi$ theory discussed above [16, 17, 24].

Focus now on BH solutions that are asymptotically flat and stationary. These admit a Killing vector ξ^μ that is timelike at infinity and acts as a generator of the event horizon. Assuming that φ respects stationarity, $\xi^\mu\nabla_\mu\varphi = 0$. Multiplying Eq. (2a) by $f_{,\varphi}$ and integrating over a volume $\mathscr{V}$ yields

$$\int_\mathscr{V} \mathrm{d}^4x\sqrt{-g}\left[f_{,\varphi}\Box\varphi + f^2_{,\varphi}(\varphi)\mathscr{G}\right] = 0\,. \tag{4}$$

Integrating by parts and using the divergence theorem, we obtain

$$\int_{\mathscr{V}} \mathrm{d}^4x\sqrt{-g}\left[f_{,\varphi\varphi}\nabla^\mu\varphi\nabla_\mu\varphi - f_{,\varphi}^2(\varphi)\mathscr{G}\right] = \int_{\partial\mathscr{V}} \mathrm{d}^3x\sqrt{|h|}f_{,\varphi}n^\mu\nabla_\mu\varphi\,, \tag{5}$$

where $\partial\mathscr{V}$ is the boundary of $\mathscr{V}$ and n^μ is the normal to the boundary. We choose $\mathscr{V}$ such that it is bounded by the BH horizon, two partial Cauchy surfaces, and spatial infinity. The contribution of the boundary term on the right-hand side vanishes. The horizon contribution vanishes by symmetry, as the normal to the horizon is ξ^μ and the stationarity condition holds; the contribution of the boundary at infinity vanishes because of asymptotic flatness. The contributions of the Cauchy surfaces exactly cancel each other, as they can be generated by an isometry. Hence the integral in the first line of Eq. (5) must vanish as well. With our signature, $\nabla^\mu\varphi\nabla_\mu\varphi$ is positive in the BH exterior. Indeed, whenever

$$f_{,\varphi\varphi}\,\mathscr{G} < 0 \tag{6}$$

the whole integrand is sign-definite and must vanish at every point in $\mathscr{V}$. The same conditions imply that the two terms of the integrand have the same sign and hence must vanish separately. This can only be achieved if $\varphi = \varphi_0$.

The above can be considered as a no-hair theorem for stationary, asymptotically flat BHs in theories that satisfy the conditions of Eqs. (3) and (6). The former is clearly an existence condition for GR solutions. To understand the latter, it is helpful to linearize Eq. (2a) around $\varphi = \varphi_0$:

$$\left[\Box + f_{,\varphi\varphi}(\varphi_0)\mathscr{G}\right]\delta\varphi = 0\,. \tag{7}$$

The term $-f_{,\varphi\varphi}\,\mathscr{G}$ acts as an effective mass $m_{\rm eff}^2$ for the perturbations $\delta\varphi$. Theories for which *this effective mass is negative* can evade the theorem above. There is a direct analogy between the proof presented here and the no-hair theorem proof of [14] for scalar-tensor theories with self-interactions.

This no-hair theorem identifies theories that can lead to interesting phenomenology in the strong-field regime: they must satisfy condition (3) but violate condition (6). A negative effective mass is expected to trigger a tachyonic instability, which can lead to the development of *scalar hair*. This is analogous to spontaneous scalarization for neutron stars (NSs) in standard scalar-tensor theories [25]. Scalarization was also shown to be possible for BHs if they are surrounded by matter [20, 21].

Quadratic scalar-Gauss-Bonnet gravity. The simplest coupling function which satisfies Eq. (3) and can violate Eq. (6) is

$$f = \eta\varphi^2/8\,, \tag{8}$$

where η is a parameter with dimensions L^2. Hereafter we will focus on this theory, and we will call it *quadratic sGB (qsGB) gravity*. If f satisfies the condition (3) and is well behaved around φ_0, then it admits the expansion $f(\varphi) = f(\varphi_0) + f_{,\varphi\varphi}(\varphi_0)(\varphi-\varphi_0)^2/2 + \ldots$ The first term in this expansion does not contribute to the field equations because $\mathscr{G}$ is a total divergence. Moreover, the kinetic term of the action is shift-symmetric. So, the field redefinition $\varphi \to \varphi - \varphi_0$ can reduce the quadratic expansion of any theory to qsGB.

qsGB gravity has several other interesting features. It leads to a field equation for φ that is linear in φ. This will be particularly convenient when studying the zero-backreaction limit below. Additionally, the theory exhibits $\varphi \to -\varphi$ symmetry. This is important in a field theory context. It prevents the term $\varphi\mathscr{G}$, which inevitably leads to BH hair [16, 17], from appearing in the action. Note also that φ does not need to play any role in late-time cosmology, hence current weak-field and gravitational wave constraints are very weak [26–29].

We focus on spherically symmetric solutions that describe either BHs or compact stars and demonstrate that spontaneous scalarization can take place. We first consider the scalar on a GR background and show that there is an instability associated to spontaneous scalarization. We then verify our results by looking at non-perturbative solutions. We call the solution with a non-trivial scalar configuration the *scalarized solution*. We focus on solutions that share the same asymptotics with the GR solution, including the asymptotic value of φ, φ_∞. For simplicity, we impose $\varphi_\infty = 0$, but this choice does not crucially affect our results.

Tachyonic instability: a zero-backreaction analysis. We first consider the limit where backreaction from the metric can be neglected, i.e. we focus on the scalar field equation, Eq. (7), on a fixed background. The effective mass of the perturbation $\delta\varphi$ is $m_{\rm eff}^2 = -f_{\varphi\varphi}\mathscr{G} = -\eta\,\mathscr{G}/4$, therefore tachyonic instability should be possible for $\eta > 0$. On a static, spherically symmetric background spacetime $\mathrm{d}s^2 = -a(r)\mathrm{d}t^2 + b(r)\mathrm{d}r^2 + r^2\mathrm{d}\Omega$, Eq. (7) can be written as

$$-\frac{\partial^2\sigma}{\partial t^2} + \frac{\partial^2\sigma}{\partial r_*^2} = V_{\rm eff}\,\sigma\,, \tag{9}$$

where $\delta\varphi = \sigma(t,r)Y_{\ell m}(\theta,\phi)/r$, $Y_{\ell m}$ are standard spherical harmonics, $\mathrm{d}r/\mathrm{d}r_* \equiv \sqrt{a/b}$ and the effective potential $V_{\rm eff}$ is:

$$V_{\rm eff} \equiv a\left[\frac{\ell(\ell+1)}{r^2} + \frac{1}{2ra}\frac{\mathrm{d}(ab^{-1})}{\mathrm{d}r} - \frac{\eta\,\mathscr{G}}{4}\right]\,. \tag{10}$$

In order to find whether scalarized solutions of the decoupled field equation (9) exist, we have performed a numerical integration, assuming a Schwarzschild background and monopolar perturbations. We have found that the equation admits a non-trivial solution with $\varphi_\infty = 0$ for a discrete spectrum of values of the coupling parameter ($\eta/M^2 = 2.902$, 19.50, $50.93, \ldots$). These results are summarized in Fig. 1, where we show the quantity $d\sigma/dr$ computed at some extraction radius $r_{\rm max} \gg M$ (namely $r_{\rm max} = 200\,M$), as a function of η/M^2. For $r \gg M$, $\delta\varphi \sim \delta\varphi_\infty + O(r^{-1})$, thus $\delta\varphi_\infty \sim d\sigma/dr(r\to\infty)$. The scalarized solutions correspond to the cusps in the top panel of Fig. 1. These solutions can be characterized by an order number $n = 0, 1, \ldots$, which is also the number of nodes of the radial profile of $\delta\varphi(r)$ (bottom-right panel of Fig. 1).

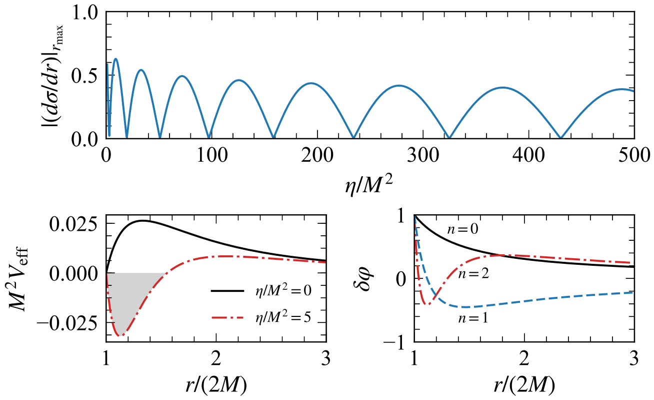

FIG. 1. *Scalar field in the decoupling limit*. Results of the numerical integration of the decoupled scalar field equation (9), assuming $\ell = 0$ and a Schwarzschild background. Top panel: asymptotic value of the scalar field as a function of η/M^2. Cusps correspond to scalarized solutions. Bottom-left panel: effective potential $V_{\rm eff}$ for $\eta/M^2 = 0$ and 5. In the latter case $V_{\rm eff}$ develops a negative region and it can support bound states. Bottom-right panel: radial profiles of $\delta\varphi$ for the first three scalarized solutions, corresponding to $\eta/M^2 = 2.902$, 19.50 and 50.93. These profiles have 0, 1 and 2 nodes, respectively.

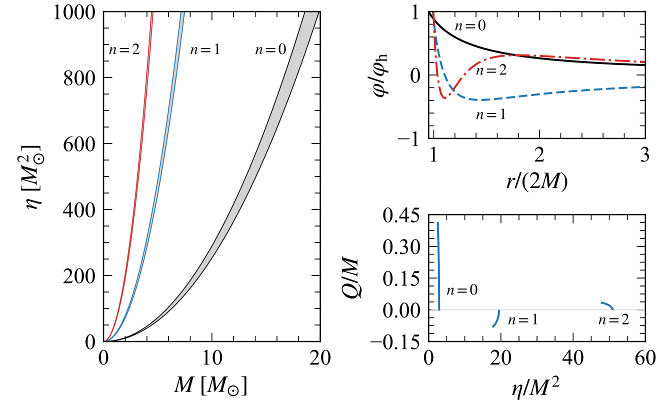

FIG. 2. *Spontaneous scalarization of black holes.* Left: the regions in the $\eta - M$ (in solar mass units) space where scalarized BHs exist. The solutions belonging to each band are characterized by the number of nodes of the scalar field radial profile. We only show the first three scalarization regions, but our numerical analysis suggests an infinite number of them. Top-right: the scalar field profiles for sample BH solutions in each of the first three bands. Bottom-right: normalized scalar charge Q/M as a function of η/M^2. The most charged BHs belong to the $n = 0$ band.

Scalarized black holes in qsGB gravity. We now consider BH solutions obtained by integrating the full set of equations (2a) and (2b). We search for static, spherically symmetric solutions, i.e. $a = a(r)$, $b = b(r)$, $\varphi = \varphi(r)$. We define $\Gamma = \log a$, $\Lambda = \log b$, as in [24]. The field equations can be cast as three coupled ordinary differential equations for Γ, Λ and φ. Since these equations are not particularly illuminating, we do not present them here.

The equation for Λ can be integrated algebraically [16, 17, 24]:

$$e^{\Lambda} = \frac{-A + \delta\sqrt{A^2 - 4B}}{2}, \quad \delta = \pm 1\,, \tag{11}$$

where $A = (1/4)r^2\varphi'^2 - (r + \eta\varphi\varphi'/2)\Gamma' - 1$ and $B = (3/2)\Gamma'\varphi'\varphi$. In BH solutions $\exp(-\Lambda)$, $\exp(\Gamma) \to \infty$ at the event horizon $r_{\rm h}$, and this implies $\delta = 1$ [24]. Replacing Eq. (11) in the remaining equations, we are left with two differential equations for Γ and φ. A near-horizon expansion of the field equations shows that $\varphi''_{\rm h} = \varphi''(r = r_{\rm h})$ is finite if

$$\varphi'_{\rm h} = \frac{r_{\rm h}}{\eta\varphi_{\rm h}}\left(-1 + \xi\sqrt{1 - 6\eta^2\varphi_{\rm h}^2/r_{\rm h}^4}\right), \tag{12}$$

where $\xi = \pm 1$. The $\xi = -1$ branch does not result in a BH solution, as discussed in [24] for the exponential coupling. Therefore, regularity on the horizon requires

$$r_{\rm h}^4 - 6\eta^2\varphi_{\rm h}^2 \geq 0\,. \tag{13}$$

Eq. (13) defines a region in the $(r_{\rm h}, \varphi_{\rm h})$–plane within which BH solutions with a regular (real) scalar field configuration exist.

The value of the scalar field at the horizon is bound in the range $0 \leq \varphi_{\rm h} \leq \varphi_{\rm h}^{\rm max} = r_{\rm h}^2/(\sqrt{6}\eta)$. We do not consider solutions with $\varphi_{\rm h} < 0$ because qsGB gravity is invariant under $\varphi \to -\varphi$. The field equations are invariant under the rescalings $r_{\rm h} \to r_{\rm h}/l$, $M \to M/l$, $\eta \to \eta/l^2$, corresponding to a freedom in choosing length units. BH solutions are then characterized by dimensionless quantities such as η/M^2 and $\eta/r_{\rm h}^2$.

For each value of η/M^2 we have numerically solved the field equations, with $\varphi_{\rm h}$ in the range $[0, \varphi^{\rm max}]$ and the other boundary conditions fixed from the requirement of regularity at the horizon. We have then extracted the scalar quantities characterizing the solution – the mass M, the scalar charge Q, and the asymptotic value of the scalar field φ_∞ – from the asymptotic expansions [17, 24, 30]:

$$e^{\Gamma} = 1 - 2M/r + Q^2M/(12r^2)\,, \tag{14}$$

$$\varphi = \varphi_0 + Q/r + QM/r^2 + (32QM^2 - Q^3)/(24r^3)\,. \tag{15}$$

While the Schwarzschild solution ($\varphi_{\rm h} = 0$, $\varphi_0 = 0$) is allowed for any value of η, a solution with $\varphi_{\rm h} \neq 0$, $\varphi_\infty = 0$ only exists when η/M^2 belongs to a set of *scalarization bands*, i.e. $[2.53, 2.89]$, $[17.86, 19.50]$, $[47.90, 50.92]$, etc. The right-end values of these bands correspond to the eigenvalues of η/M^2 found by solving the linear equation of the scalar field on a fixed background. The scalarization bands in η/M^2 correspond to regions bounded by parabolas in the $\eta - M$ plane (shadowed regions in the left panel of Fig. 2). The scalar field profiles of these solutions have $n = 0, 1, \ldots$ nodes (top-right panel of Fig. 2), corresponding to the order number of the scalarization band. A similar ladder of excited states was observed for scalarized NSs in scalar-tensor theory [31, 32]. The normalized scalar charge[1] Q/M of these solutions is shown in

[1] In other theories with a Gauss–Bonnet coupling the scalar charge and the

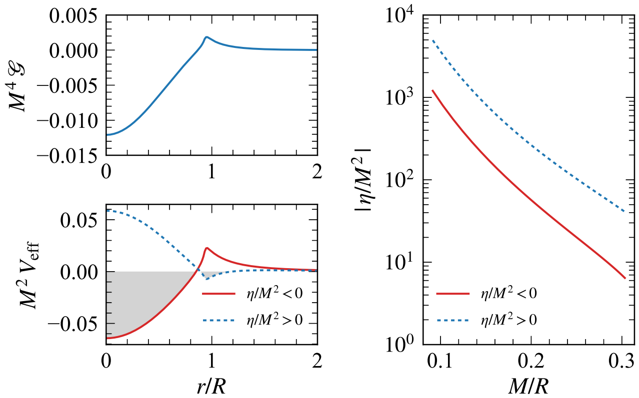

FIG. 3. *Tachyonic instability in a stellar background.* Left: profile of the Gauss-Bonnet invariant (top) and of the effective potential (bottom), inside a $M = 1.4\ \mathrm{M}_\odot$ NS with the SLy4 EoS, assuming $|\eta/M^2| = 100/(1.4)^2 \sim 51$. The regions where the effective potential becomes negative are shaded. Right: values of η/M^2 for which the first bound state forms as a function of the compactness M/R.

the bottom-right panel of Fig. 2 as a function of η/M^2. This plot shows the values of η admitting a scalarized solution for each value of the BH mass.

Spontaneous scalarization and neutron stars. Let us now consider NSs in qsGB gravity. The Gauss-Bonnet invariant for a static, spherically symmetric solution of the Tolman-Oppenheimer-Volkoff (TOV) equations [33] is

$$\mathscr{G} = \frac{48m^2}{r^6} - \frac{128\pi(m + 2\pi r^3 p)\varepsilon}{r^3}, \tag{16}$$

where $m = r(1 - 1/b)/2$ is the mass function, and p and ε are the pressure and energy density inside the star, respectively. At the surface $r = R$, ε vanishes and (16) matches smoothly the Schwarzschild value $\mathscr{G} = 48M^2/r^6$, with $M \equiv m(R)$ being the star's mass. We solve the TOV equations for a "canonical" NS model with $M = 1.4\ \mathrm{M}_\odot$, assuming the SLy4 [34] equation of state (EoS). The Gauss-Bonnet invariant is mostly negative throughout the interior of the star (see Fig. 3, top-left panel); it is only positive near the surface of the star, and in the exterior. This suggests that if $\eta < 0$, the scalar field can develop a tachyonic instability inside the star, while if $\eta > 0$ the instability is triggered in the outer region/exterior of the star.

In the bottom-left panel of Fig. 3 we show the effective potential $V_{\rm eff}$ for the "canonical" NS model discussed above, with $\eta = \pm 100\ \mathrm{M}_\odot^2$. As expected, there are (shaded) regions where $V_{\rm eff}$ becomes negative. These regions are inside the star when $\eta < 0$, and outside the star when $\eta > 0$.

Solving Eq. (9) in the NS background, we find that scalarized solutions exist for both positive and negative values of η. In the right panel of Fig. 3 we show the values of η/M^2 corresponding to the lowest-lying scalarized solutions with $\eta > 0$ and $\eta < 0$, as a function of the NS compactness. Note that scalarization occurs for lower values of $|\eta/M^2|$ when the coupling constant is negative than when it is positive.

As in the BH case, we expect these results to translate into the existence of scalarized NSs at the fully nonlinear level [35], i.e. by integrating the modified TOV equations obtained from Eqs. (2a)-(2b) assuming a perfect fluid for matter. Fully nonlinear stellar models will be explored in forthcoming work.

Conclusions. We have identified and studied a subclass of scalar-tensor theories with a coupling between the scalar and the Gauss–Bonnet invariant that appears to exhibit spontaneous scalarization for both BHs and NSs. Interestingly, BH scalarization does not have a single threshold. Instead, for a given value of the coupling parameter η hairy BHs exist when their mass lies in one of many narrow bands. Our exploration for NSs strongly suggests that scalarization can take place for both positive and negative values of η. However, the effect appears to be stronger for negative values of η, for which BH scalarization cannot occur. A full numerical study of NSs in these theories is in progress and will be reported elsewhere. It would be interesting to examine more closely the conditions under which spontaneous scalarization can occur and its implications for the structure of astrophysical BHs and compact stars, especially in binary systems of interest for gravitational wave detectors. A full study of the two-body problem in qsGB is beyond the scope of this paper, but we anticipate interesting phenomenology **already at the post-Newtonian level [36].** Binary systems containing scalarized BHs and NSs (which have nonzero scalar charge Q) should emit dipolar scalar radiation. However, in contrast with dilatonic and shift-symmetric theories, where $Q \neq 0$ for *all* BHs, in our case scalarization only happens – and therefore dipolar radiation would be emitted – only in certain BH mass ranges (for a fixed coupling η). NSs in the shift-symmetric theory have $Q = 0$ [37, 38], thus evading the stringent experimental constraints on dipolar radiation emission **from binary pulsars** [39]. In qsGB gravity, if one of the NSs in the binary happens to be scalarized, scalar radiation would be emitted, leaving a smoking gun of the presence of the scalar field in the orbital dynamics. **It would also be interesting to investigate the strong field dynamics of this theory. Apart from scalar-tensor theories [40–43], the application of numerical relativity simulations to other theories of gravity is still in its infancy [44–47]. To perform numerical simulations one must inevitably address the issue of well-posedness [48, 49], which remains an open problem beyond the scope of our paper. By pointing out the** *existence* **of potentially interesting phenomenology in qsGB we hope to motivate further work in this direction.** Finally, it might also be worth extending our results to more general couplings between the scalar field and the Gauss–Bonnet invariant.

Note. During the completion of this manuscript, a preprint studying a similar model for BHs appeared on the arXiv [50].

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asymptotic value of the coupling are related by $Q/M = 2f_{,\varphi}(\varphi_\infty)/M^2$, and this can lead to a bound on the coupling constant (e.g. [16, 24, 30]). It should be noted that there is no such relation for qsGB because $f_{,\varphi}(\varphi_\infty) = 0$.

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[1] B. P. Abbott *et al.* (Virgo, LIGO Scientific), Phys. Rev. Lett. **116**, 061102 (2016), arXiv:1602.03837 [gr-qc].
[2] B. P. Abbott *et al.* (Virgo, LIGO Scientific), Phys. Rev. Lett. **116**, 241103 (2016), arXiv:1606.04855 [gr-qc].
[3] B. P. Abbott *et al.* (Virgo, LIGO Scientific), Phys. Rev. **X6**, 041015 (2016), arXiv:1606.04856 [gr-qc].
[4] B. P. Abbott *et al.* (VIRGO, LIGO Scientific), Phys. Rev. Lett. **118**, 221101 (2017), arXiv:1706.01812 [gr-qc].
[5] B. P. Abbott *et al.* (Virgo, LIGO Scientific), Phys. Rev. Lett. **119**, 141101 (2017), arXiv:1709.09660 [gr-qc].
[6] B. P. Abbott *et al.* (Virgo, LIGO Scientific), Astrophys. J. **851**, L35 (2017), arXiv:1711.05578 [astro-ph.HE].
[7] B. Abbott *et al.* (Virgo, LIGO Scientific), Phys. Rev. Lett. **119**, 161101 (2017), arXiv:1710.05832 [gr-qc].
[8] E. Berti *et al.*, Class. Quant. Grav. **32**, 243001 (2015), arXiv:1501.07274 [gr-qc].
[9] T. P. Sotiriou, *Proceedings of the 7th Aegean Summer School: Beyond Einstein's theory of gravity. Modifications of Einstein's Theory of Gravity at Large Distances. Paros, Greece, September 23-28, 2013*, Lect. Notes Phys. **892**, 3 (2015), arXiv:1404.2955 [gr-qc].
[10] T. P. Sotiriou, Class. Quant. Grav. **32**, 214002 (2015), arXiv:1505.00248 [gr-qc].
[11] C. A. R. Herdeiro and E. Radu, *Proceedings, 7th Black Holes Workshop 2014: Aveiro, Portugal, December 18-19, 2014*, Int. J. Mod. Phys. **D24**, 1542014 (2015), arXiv:1504.08209 [gr-qc].
[12] S. W. Hawking, Commun. Math. Phys. **25**, 167 (1972).
[13] J. D. Bekenstein, Phys. Rev. **D5**, 1239 (1972).
[14] T. P. Sotiriou and V. Faraoni, Phys. Rev. Lett. **108**, 081103 (2012), arXiv:1109.6324 [gr-qc].
[15] L. Hui and A. Nicolis, Phys. Rev. Lett. **110**, 241104 (2013), arXiv:1202.1296 [hep-th].
[16] T. P. Sotiriou and S.-Y. Zhou, Phys. Rev. Lett. **112**, 251102 (2014), arXiv:1312.3622 [gr-qc].
[17] T. P. Sotiriou and S.-Y. Zhou, Phys. Rev. **D90**, 124063 (2014), arXiv:1408.1698 [gr-qc].
[18] E. Barausse and T. P. Sotiriou, Phys. Rev. Lett. **101**, 099001 (2008), arXiv:0803.3433 [gr-qc].
[19] A. Arvanitaki, S. Dimopoulos, S. Dubovsky, N. Kaloper, and J. March-Russell, Phys. Rev. **D81**, 123530 (2010), arXiv:0905.4720 [hep-th].
[20] V. Cardoso, I. P. Carucci, P. Pani, and T. P. Sotiriou, Phys. Rev. Lett. **111**, 111101 (2013), arXiv:1308.6587 [gr-qc].
[21] V. Cardoso, I. P. Carucci, P. Pani, and T. P. Sotiriou, Phys. Rev. **D88**, 044056 (2013), arXiv:1305.6936 [gr-qc].
[22] E. Babichev and C. Charmousis, JHEP **08**, 106 (2014), arXiv:1312.3204 [gr-qc].
[23] C. A. R. Herdeiro and E. Radu, Phys. Rev. Lett. **112**, 221101 (2014), arXiv:1403.2757 [gr-qc].
[24] P. Kanti, N. E. Mavromatos, J. Rizos, K. Tamvakis, and E. Winstanley, Phys. Rev. **D54**, 5049 (1996), arXiv:hep-th/9511071 [hep-th].
[25] T. Damour and G. Esposito-Farèse, Phys.Rev.Lett. **70**, 2220 (1993).
[26] G. Esposito-Farese, *Phi in the sky: The quest for cosmological scalar fields. Proceedings, Workshop, Porto, Portugal, July 8-10, 2004*, AIP Conf. Proc. **736**, 35 (2004), [,35(2004)], arXiv:gr-qc/0409081 [gr-qc].
[27] T. P. Sotiriou and E. Barausse, Phys. Rev. **D75**, 084007 (2007), arXiv:gr-qc/0612065 [gr-qc].
[28] J. Sakstein, B. Jain, J. S. Heyl, and L. Hui, Astrophys. J. **844**, L14 (2017), arXiv:1704.02425 [astro-ph.CO].
[29] J. Sakstein and B. Jain, (2017), arXiv:1710.05893 [astro-ph.CO].
[30] S. Mignemi and N. R. Stewart, Phys. Rev. **D47**, 5259 (1993), arXiv:hep-th/9212146 [hep-th].
[31] W. C. C. Lima, G. E. A. Matsas, and D. A. T. Vanzella, Phys. Rev. Lett. **105**, 151102 (2010), arXiv:1009.1771 [gr-qc].
[32] P. Pani, V. Cardoso, E. Berti, J. Read, and M. Salgado, Phys. Rev. **D83**, 081501 (2011), arXiv:1012.1343 [gr-qc].
[33] B. K. Harrison, K. S. Thorne, M. Wakano, and J. A. Wheeler, *Gravitation Theory and Gravitational Collapse, Chicago: University of Chicago Press* (1965).
[34] F. Douchin and P. Haensel, Astron. Astrophys. **380**, 151 (2001), arXiv:astro-ph/0111092 [astro-ph].
[35] P. Pani, E. Berti, V. Cardoso, and J. Read, Phys. Rev. **D84**, 104035 (2011), arXiv:1109.0928 [gr-qc].
[36] K. Yagi, L. C. Stein, N. Yunes, and T. Tanaka, Phys. Rev. **D85**, 064022 (2012), [Erratum: Phys. Rev.D93,no.2,029902(2016)], arXiv:1110.5950 [gr-qc].
[37] E. Barausse and K. Yagi, Phys. Rev. Lett. **115**, 211105 (2015), arXiv:1509.04539 [gr-qc].
[38] K. Yagi, L. C. Stein, and N. Yunes, Phys. Rev. **D93**, 024010 (2016), arXiv:1510.02152 [gr-qc].
[39] P. C. Freire, N. Wex, G. Esposito-Farèse, J. P. Verbiest, M. Bailes, *et al.*, Mon. Not. Roy. Astron. Soc. **423**, 3328 (2012), arXiv:1205.1450 [astro-ph.GA].
[40] E. Barausse, C. Palenzuela, M. Ponce, and L. Lehner, Phys. Rev. **D87**, 081506 (2013), arXiv:1212.5053 [gr-qc].
[41] E. Berti, V. Cardoso, L. Gualtieri, M. Horbatsch, and U. Sperhake, Phys. Rev. **D87**, 124020 (2013), arXiv:1304.2836 [gr-qc].
[42] M. Shibata, K. Taniguchi, H. Okawa, and A. Buonanno, Phys. Rev. **D89**, 084005 (2014), arXiv:1310.0627 [gr-qc].
[43] C. Palenzuela, E. Barausse, M. Ponce, and L. Lehner, Phys. Rev. **D89**, 044024 (2014), arXiv:1310.4481 [gr-qc].
[44] R. Benkel, T. P. Sotiriou, and H. Witek, Phys. Rev. **D94**, 121503 (2016), arXiv:1612.08184 [gr-qc].
[45] R. Benkel, T. P. Sotiriou, and H. Witek, Class. Quant. Grav. **34**, 064001 (2017), arXiv:1610.09168 [gr-qc].
[46] M. Okounkova, L. C. Stein, M. A. Scheel, and D. A. Hemberger, Phys. Rev. **D96**, 044020 (2017), arXiv:1705.07924 [gr-qc].

[47] E. W. Hirschmann, L. Lehner, S. L. Liebling, and C. Palenzuela, (2017), arXiv:1706.09875 [gr-qc].

[48] G. Papallo and H. S. Reall, Phys. Rev. **D96**, 044019 (2017), arXiv:1705.04370 [gr-qc].

[49] J. Cayuso, N. Ortiz, and L. Lehner, Phys. Rev. **D96**, 084043 (2017), arXiv:1706.07421 [gr-qc].

[50] D. D. Doneva and S. S. Yazadjiev, (2017), arXiv:1711.01187 [gr-qc].