Gravitational Waves from Accretion-Induced Descalarization in Massive Scalar-Tensor Theory

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Many classes of extended scalar-tensor theories predict that dynamical instabilities can take place at high energies, leading to the formation of scalarized neutron stars. Depending on the theory parameters, stars in a scalarized state can form a solution-space branch that shares a lot of similarities with the so-called mass twins in general relativity appearing for equations of state containing first-order phase transitions. Members of this scalarized branch have a lower maximum mass and central energy density compared to Einstein ones. In such cases, a scalarized star could potentially overaccrete beyond the critical mass limit, thus triggering a gravitational phase transition where the star sheds its scalar hair and migrates over to its nonscalarized counterpart. Such an event resembles, but is distinct from, a nuclear or thermodynamic phase transition. We dynamically track a gravitational transition by first constructing hydrostatic, scalarized equilibria for realistic equations of state, and then allowing additional material to fall onto the stellar surface. The resulting bursts of monopolar radiation are dispersively stretched to form a quasicontinuous signal that persists for decades, carrying strains of order $\gtrsim 10^{-22}\ (\mathrm{kpc}/L)^{3/2}\,\mathrm{Hz}^{-1/2}$ at frequencies of $\lesssim 300$ Hz, detectable with the existing interferometer network out to distances of $L \lesssim 10$ kpc, and out to a few hundred kpc with the inclusion of the Einstein Telescope. Electromagnetic signatures of such events, involving gamma-ray and neutrino bursts, are also considered.

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Introduction.—General relativity (GR) has historically provided an excellent description for both local (e.g., Solar System) and global (e.g., cosmology) gravitational phenomena. It is well known, however, that the theory cannot by itself be fully complete, and the nonrenormalizability of the action implies that additional ingredients, possibly in the form of nonminimally coupled fields [1,2], should activate at extreme scales. Nevertheless, any theoretical extension must be virtually invisible at low energies, and also somehow suppressed in certain strong-field environments. For example, binary pulsar experiments restrict the possibility for significant subquadrupolar radiation over super-Compton length scales [3–7], and gravitational-wave (GW) experiments suggest that (at least some) black holes should be approximately, if not exactly, Kerr [8]. An observationally viable class of extensions that can survive these issues is massive scalar-tensor theory (STT): the mass of the scalar field suppresses the scalar dipole radiation [9,10], and the classical no-hair theorems tend to be respected [11], implying that astrophysically stable black holes would be indistinguishable from their GR counterparts.

Material degrees of freedom in these theories, however, allow for the possibility of "scalarized" stars [12]. This phenomenon can be generally thought of as a consequence of the effective curvature-coupled mass term, which appears in the relevant Klein-Gordon equation, changing sign once some critical threshold is breached, thereby inducing a tachyonic instability (also see, though, Ref. [13]). Thus, for neutron stars a critical compactness exists at which a branch of strongly scalarized solutions emerges. In some cases, the heaviest scalarized neutron star has lower baryon mass and central energy density compared to the maximum mass (stable) nonscalarized neutron star [14] and there is a gap between the two branches where no stable neutron star solutions exist. This picture resembles very closely the so-called mass twins in pure GR that are manifestations of the presence of a first-order phase transition in the equation of state (EOS) [15–18]. This implies that if a near-critical scalarized star were to acquire additional mass through accretion, the system may promptly discharge its scalar hair. Importantly, the neutron star does not need to collapse to a black hole in this scenario, as considered in, e.g., [19–22], but rather may undergo a "gravitational" phase transition, distinct from a "material" (e.g., hadron-quark) phase transition [23], and migrate to the GR branch pertaining to the same EOS. This novel scenario is considered in this Letter.

A migration to a new branch is likely to carry a variety of observational signatures. As a scalar shedding necessarily compactifies the star over a short, dynamical timescale ($\gtrsim$ms), abrupt changes in the electromagnetic output of the source, most notably associated with gamma-ray burst (GRB) afterglows and neutrino bursts [24–26], would point toward such a transition. These signatures could, however, be imitated by a nuclear phase transition [27–32], thus highlighting the well-known degeneracy between effects coming from a modification of gravity and the uncertainties in the nuclear EOS (e.g., [33,34]). One key difference is that a gravitational transition will unleash a burst of scalar radiation which, for a massive scalar field theory, will be dispersively stretched into a quasicontinuous signal, as higher frequency components are first to arrive at the detector(s) [21]. GW afterglows lasting up to a ~kyr may therefore follow a gravitational phase transition.

Formalism and equations of motion.—The action of a STT of whichever flavor [e.g., Brans-Dicke, Bergmann-Wagoner, or even $f(R)$ theories] can be transformed into the Einstein frame

$$S = \int \frac{\sqrt{-g}}{16\pi} d^4x(R - 2\partial_\mu\varphi\partial^\mu\varphi - 4V) + S_M[\Psi, A^2 g_{\mu\nu}] \quad (1)$$

for matter portion S_M, metric tensor g, scalar field φ, Ricci scalar R. Scalar field potential is taken as $V(\varphi) = m_\varphi^2\varphi^2/2$ [35], whose saddle point poses the boundary condition $\varphi(r \to \infty) = 0$ for φ. The transition to the physical Jordan frame metric $\tilde{g}_{\mu\nu}$ is done via a Weyl scaling $\tilde{g}_{\mu\nu} = A(\varphi)^2 g_{\mu\nu}$, where we adopt the spherically symmetric Jordan frame metric, $\tilde{g}_{\mu\nu} = \mathrm{diag}[-\alpha^2, X^2, A(\varphi)^2 r^2, A(\varphi)^2 r^2 \sin^2\theta]$, and choose the conformal factor $A(\varphi) = \exp(\alpha_0\varphi + \beta_0\varphi^2/2)$ following [35,36]. The resulting field equations are given in the Supplemental Material [37], while we note here that the term responsible for scalarization is proportional to the logarithmic derivative of the conformal factor $\alpha(\varphi) = d\ln A/d\varphi$. The choice we make, $\alpha(\varphi) = \alpha_0 + \beta_0\varphi$, therefore corresponds to the two leading terms in the Maclaurin expansion of any regular function; higher-order terms do not change the picture of scalarization qualitatively for a large class of more complicated conformal factors because $\varphi \ll 1$ everywhere [12,48,49]. Simulations for another coupling function are shown in the Supplemental Material [37] to demonstrate that the phenomenon put forward in this Letter is not quenched by higher-order effects.

Two additional GW modes beyond those in GR are raised by φ, viz. the breathing ("B") and longitude ("ℓ") modes, which carry the same response functions up to a sign flip [see Eq. (134) of [50]]. The strain of the latter is weaker by a factor $(\lambda_\varphi f)^{-2}$ relative to the former, which reads $h_\varphi = 2\alpha_0\varphi$. The strain felt by a LIGO-Virgo-like array (two orthogonal antennas), at a distance L from the source, is thus

$$h(L,t) = h_\varphi(L,t)\{1 - [\lambda_\varphi f(L,t)]^{-2}\}/2 \quad (2)$$

when the source orients optimally, where $\lambda_\varphi = 2\pi\hbar/m_\varphi$ is the Compton length scale for the massive scalar field, t is the retarded time postemission, and $f(L,t)$ is the characteristic frequency of the signal. Furthermore, modes with distinct frequencies propagate at different, subluminal velocities, and the full power-spectral density (PSD), $2\sqrt{f}|\tilde{h}(f)|$, will not arrive at the detector simultaneously. As a result, the dispersively stretched burst becomes quasimonochromatic over a L- and m_φ-dependent timescale [21,35]. Therefore, there is an implicit time dependence, encoded in f, for the witness. The quasimonochromatic feature implies that a phase-coherent search can be implemented, and the signal-to-noise ratio (SNR), $4\int df[|\tilde{h}(f)|^2/S_n(f)]$, can be obtained by integrating the strain over a narrow frequency interval. In the limit $Tf \gg 1$, i.e., when many cycles are observed, the SNR can be obtained by dividing the "effective" PSD from the aforementioned time-domain integration [21],

$$\sqrt{S_f} = \sqrt{T}\alpha_0 A(L,t)\{1 - [\lambda_\varphi f(L,t)]^{-2}\}, \quad (3)$$

by the noise spectral curve $\sqrt{S_n(f)}$, where we assume an optimally oriented detector. Here, the amplitude is $A(L,t) \simeq \mathcal{A}(f)L^{-3/2}\lambda_\varphi(f^2 - \lambda_\varphi^{-2})^{3/4}$ with $\mathcal{A}(f)$ the Fourier component of the scalar GW extracted at some distance $\lambda_\varphi < r_{\mathrm{out}} \ll L$ so that it contains the wave content that eventually propagates to a detector at L [cf. Eq. (57) in [21], while noting that our definition for A differs from theirs by a factor of L^{-1}].

Scalarized neutron stars.—In the present Letter, we adopt a piecewise-polytropic approximation to the APR4 EOS [51], which withstands constraints coming from GW 170817 [52] and supports masses that accommodate the heaviest neutron star observed to date, viz. PSR J0740 + 6620 ($M = 2.14^{+0.10}_{-0.09}$ $M_\odot$) [53]. Specifics related to this EOS and its numerical implementation are given in the Supplemental Material [37]. In the considered STT and for large enough stellar compactness, the neutron star scalarizes, meaning it develops a strong, localized scalar field. If $\alpha_0 \neq 0$, purely GR solutions do not exist within the theory, and all stars have at least some tiny residual $\varphi \neq 0$, i.e., they must be at least weakly scalarized. For practical purposes, however, weakly scalarized solutions are virtually indistinguishable from GR counterparts at high densities and thus, with a slight abuse of language, we call these solutions "descalarized."

A sequence of neutron star solutions is shown in the top panel of Fig. 1 as a function of central energy density ϵ_c for

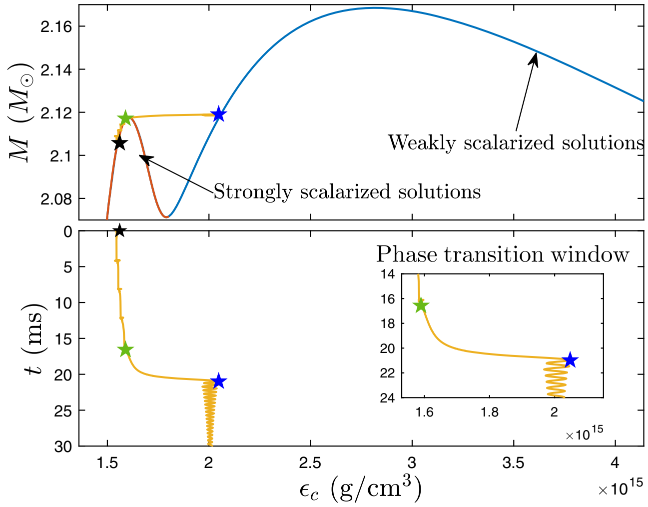

FIG. 1. Evolutionary track of an APR4, near-critical scalarized star under accretion: gravitational mass, M, as a function of central energy density, ϵ_c (top panel), together with the time evolution of ϵ_c itself (bottom panel). The blue branch represents weakly scalarized stars, while the red branch is strongly scalarized. The green and blue stars mark the onset and the termination of descalarization, respectively, while the black star marks the initial state of the accretion simulation described in the main text.

$\alpha_0 = 10^{-2}$, $\beta_0 = -5$, and $m_\varphi = 10^{-14}$ eV. This value of m_φ mitigates the tension with binary-pulsar constraints, since radiation is suppressed over super-Compton length scales $r \gg \lambda_\varphi$ [4,9,10], and effectively allows for a broad range of β_0 [54]. The other parameters, which similarly respect binary-pulsar constraints, are chosen such that the maximal mass for the scalarized branch (2.118 $M_\odot$; red curve) is less than that of the "weakly scalarized" branch (2.168 $M_\odot$; blue curve). Moreover, for both the red and the blue branches the neutron star solutions are stable up to the maximum mass point and lose stability afterward [22,55,56], resulting in a picture reminiscent of the so-called mass twins in pure GR [15–18]. Thus, a phase transition (descalarization) from the red to the blue branch can be realized if additional mass is added. For the APR4 EOS, this scenario is possible provided that $-5.41 \lesssim \beta_0 \lesssim -4.78$, with the exact range depending on α_0 and the coupling functions $V(\varphi)$ and $A(\varphi)$. Similar ranges apply for other EOSs. These limits will differ also if we consider scalarization in more general theories of gravity, such as tensor-multiscalar theories [57–59].

Accretion dynamics.—A neutron star that overaccretes beyond the peak of the scalarized curve, displayed by the green star in Fig. 1, will undergo a phase transition by descalarizing. This scenario may occur either for a newborn star after a merger or collapse through fallback accretion, or a mature star in a binary undergoing Roche-lobe overflow. In the former case, debris disks containing $\lesssim 0.2\ M_\odot$ worth of material [25], though potentially much more in a core collapse [60], will form around the birth site. A sizeable fraction ($\lesssim 40\%$) of it may eventually fall back onto the stellar surface [61]. Accreted masses may total $\lesssim 0.8\ M_\odot$ in some x-ray binaries [62], though such amounts can accumulate only over long (potentially ~Gyr) timescales. The details of the accretion process itself are complicated, however, since the neutron star may be spinning rapidly enough that material is repelled by a centrifugal barrier ("propeller" effect [60]), pressure gradients from nucleosynthetic heating can accelerate ejecta before it has a chance to return [63], and material will not fall isotropically onto the surface but rather may be guided onto the magnetic poles by the Lorentz force [64].

In this Letter, however, our main goal is not to simulate a realistic accretion process in a STT, but rather to illustrate qualitatively how the dynamical acquisition of additional mass can trigger a descalarization. To this end, accretion is artificially simulated by superposing a radial, Gaussian bulk centred at 0.9 $R_\star$, for stellar radius $R_\star$, with a width ("standard deviation") of 1 km, every 4 ms. The process is then halted when a total (baryon) mass of 0.015 $M_\odot$ has been added (after 16.01 ms). The average accretion rate of $\simeq 0.94\ M_\odot\,s^{-1}$ is marginally slower than that observed in the first few ms of merger simulations, viz. $\lesssim 1\ M_\odot\,s^{-1}$ (Fig. 7 of [65]). Using this scheme, we model a dynamical descalarization. As shown in Fig. 1, the system begins in a particular state (shown by the black star), gains some mass (green star), and the descalarization process begins (orange line), until eventually the system oscillates around a certain, stable state on the new branch (blue star). The bottom panel of Fig. 1 shows the evolution of ϵ_c in this example; the increase of ϵ_c from 1.6×10^{15} g cm^{-3} to $\sim 2\times10^{15}$ g cm^{-3} in ~ 4 ms indicates that a rapid compactification accompanies descalarization.

We have verified that the (post)descalarization dynamics remain the same if the bulk is accreted with longer waiting time, and that our results are not overly sensitive to the particulars of the chosen accretion profile, as described above, by studying other bulk impositions (see the Supplemental Material [37]). Nevertheless, we stress that our profiles are not representative of realistic astrophysical processes, though they allow us to capture the salient features of a gravitational phase transition. Importantly, magnetic fields only couple weakly to the scalar sector, and thus even if the geometry of the accreted-mass buildup ("mountain") is sensitive to the former (e.g., [66]), the descalarization dynamics, and resulting scalar-GW signal, are not. The latter aspects are discussed below.

Results.—Introducing the auxiliary variables $\psi = \alpha^{-1}(\partial\varphi/\partial t)$ and $\eta = X^{-1}(\partial\varphi/\partial r)$, we have that the energy E_φ and luminosity $\mathcal{L}_\varphi$ of the scalar field read (see Supplemental Material [37])

$$E_\varphi = \int dr\left[\frac{r^2 A(\varphi)^2}{2}(\psi^2+\eta^2)+V\right] \quad (4)$$

and

$$\mathcal{L}_\varphi = A(\varphi)^{-2} r^2 X \alpha \psi \eta, \tag{5}$$

which defines the corresponding energy leakage as

$$E_{\rm GW}^{\rm (scalar)} = \int \mathcal{L}_\varphi dt. \tag{6}$$

For the simulation shown in the top panel of Fig. 1, the scalar energy [Eq. (4)] plummets to zero, from its initial value of 0.051 $M_\odot$, after descalarization. The near zone ($r \ll \lambda_\varphi$) extraction of $E_{\rm GW}^{\rm (scalar)}$ suggests an energy loss $\gtrsim 40$ times less than the decrease of E_φ, indicating that most of the scalar energy transforms into gravitational binding energy since the stellar radius shrinks from 11.56 to 10.36 km between the initial and final states, while the (gravitational) mass increases by ≈ 0.013 $M_\odot$. After the emission propagates to distances comparable to λ_φ, the dispersion suppresses the low frequency component(s) and leads to a stretching of the waveform. The scalar energy leakage ultimately saturates at super-Compton length scales, where a reduction of factor ~ 2 in the near-zone radiation is seen due to the dispersion.

Assuming an observation duration of $T = 2$ months, over which the signal evolves slowly for $m_\varphi = 10^{-14}$ eV [21], we present the numerical evaluation of the signal amplitude [Eq. (3)] for scalar GWs for $\alpha_0 = 10^{-2}$ at a fixed distance $L = 10$ kpc in Fig. 2. The kth notch, starting from the right, plotted over the signal (black curve) stands for $t = 10^{(k-1)}$ years, illustrating that the bulk of the signal persists for $\lesssim$ centuries. We see that for $T = 2$ months and these theory parameters (see also below), the signal should be detectable with sufficiently high signal-to-noise ratio with the existing interferometer network out to distances of $L \lesssim 10$ kpc, and out to a few hundred kpc with the inclusion of the Einstein Telescope (ET). Note that the distance and observation time T are, to a large degree, degenerate. We find that increasing T by factor ~ 2 allows for events at distances of factor $\sim 2^{1/3}$ farther out to become visible, assuming a narrow-band search is carried out.

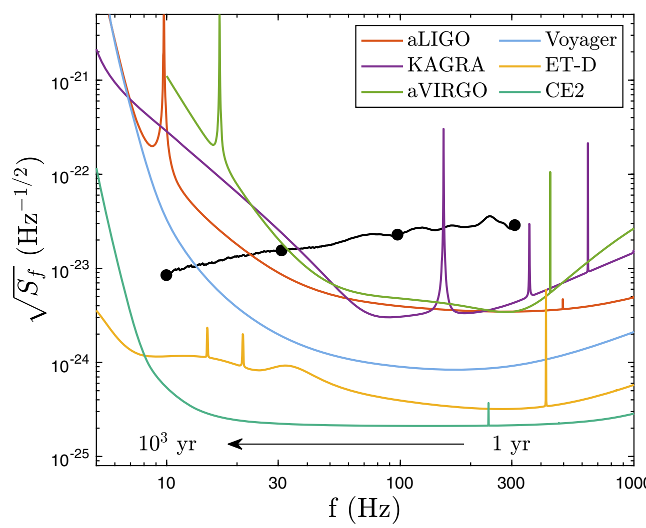

FIG. 2. Effective, root PSD $\sqrt{S_f}$ [Eq. (3)] at $L = 10$ kpc and $T = 2$ months for the strain of the scalar-induced GW mode as a function of retarded time from 1 to 10^3 years (black curve; the kth dot from the right along the curve represents $10^{(k-1)}$ years). Overlaid are the sensitivity curves of existing and upcoming GW interferometers.

To quantify the detectability in general, we compared outputs from a variety of simulations with varying α_0. We find a fitting to the root of the "effective" PSD, in units of Hz$^{-1/2}$, as

$$\frac{\sqrt{S_f}}{10^{-23}} \approx 3 \left(\frac{\alpha_0}{10^{-2}}\right) \left(\frac{T}{5 \times 10^6\ \rm s}\right)^{\frac{1}{2}} \left(\frac{10\ \rm kpc}{L}\right)^{\frac{3}{2}} \left(\frac{f}{300\ \rm Hz}\right)^{0.34}, \tag{7}$$

with the frequency of the signal approximated as $f(L,t) \approx 2.42(t+L)/\sqrt{t(t+2L)}$ Hz [cf. Eq. (53) in [21]]. We note that larger T may be used for greater retarded times since the timescale for the frequency evolution, $f/\dot{f}$, scales as t, thus opening the possibility for much larger effective PSD. In particular, since the signal [Eq. (7)] is quasicontinuous, an extended narrow-band search could be carried out if one knew when the system descalarized, as the dispersion relation directly equates the relative delay with a frequency. Furthermore, multiple sensors can act to "fuse" data together in a way that improves the overall signal-to-noise ratio beyond that inferred from Eq. (7) (see Sec. IV. E. of [67] for a detailed comparison of achievable sensitivities with different networks).

Connection to matter phase transitions.—The phase transitions from scalarized to nonscalarized states considered here bear striking similarities to the material phase transitions from confined hadronic to deconfined quark matter. In both cases, there can be stars of equal mass but different radii that are separated by a range of central energy densities where the stable solution space is empty (i.e., twin stars [15–18]). The astrophysical implications of matter phase transitions, and especially the GW signatures, have attracted considerable attention recently (see, e.g., Refs. [27–32]). In each case there will be a descalarization analog, with the main difference being an additional channel for energy loss: the scalar radiation.

As a proof of principle we concentrate on accretion in this Letter. However, such analogs can be found also in cases without accretion. For a hot, newborn neutron star with an EOS that permits negatively charged, nonleptonic particles (e.g., hyperons or quarks), the hydrostatic support available to the star will reduce when neutrinos diffuse out of the core [68]. This can lead to a delayed phase transition with a number of interesting observational signatures [29].

A descalarization analog of this delayed transition exists: depending on the chemical composition and theory parameters, the scalarized star may migrate to a non- or weakly scalarized branch when the temperature drops below a critical threshold. A similar picture exists for cases where the star is centrifugally or magnetically supported: spin-down or field decay reduces the maximum mass of the system, which could force the star to transition [69,70]. Studying these processes in detail lies beyond the scope of the present Letter, though complementing scalar-flavor phase transitions with studies of neutron star mergers in STT is likely to offer rich phenomenology as concerns the evolutionary track of neutron stars. This will be, on one hand, due to the additional channel of energy loss that, even if not detectable, will alter the merger remnant evolution. Some properties of the postmerger remnant, such as its oscillations frequencies, can also differ from GR due to the scalarization-related changes in stellar structure [71].

We point out that we discuss twin stars only as an interesting analogy with the observed process of descalarization. Our simulations and the predicted observational signatures are completely independent of the existence of such stars (the astrophysical relevance of twin stars is discussed in, e.g., [72–75]).

Discussion and observational prospects.—While a detection of scalar GWs of the form shown in Fig. 2 could be used to unambiguously identify that a descalarization took place, (massive) STTs may already leave traces in the events that lead up to the transition. A promising avenue for the formation of scalarized stars, which are also prone to overaccretion and descalarization, comes from binary mergers. The scalar field associated with the binary constituents may become excited during inspiral, leaving a clear imprint on the GW signal by accelerating the coalescence [71] (see also Refs. [76–80]). Despite progress, though, certain key effects such as rotation (see, however, Refs. [81]) are still missing from numerical simulations of mergers involving stars in STTs. This means that direct waveform comparisons with observed inspirals cannot be achieved yet. On the electromagnetic side, however, binary neutron-star merger events are also the progenitors for short GRBs, which offer avenues for indirectly observing a descalarization.

Many GRBs exhibit extended emissions at short wavelengths following the main burst. Emission profiles that display a long-lived x-ray "plateau" are suggestive of persistent energy injections ("magnetar wind") from a massive, newborn neutron star [24,25]. Suppose that tensorial GWs were coincidentally observed with a short GRB (as occurred for GW170817 [82]), followed by a plateaulike x-ray afterglow. The detection of a "scalar" GW afterglow some time after the main event, which may persist for $\lesssim$ centuries (see Fig. 2), would clearly indicate that the remnant peeled its scalar hair. Even without such measurements, the nature of the electromagnetic afterglow will be affected by a scalar shedding as the star condenses (see Fig. 1). The spin-down power associated with magnetic dipole braking scales as $L_{\rm dip} \propto R_\star^6$ (e.g., [25]), and so a decrease in $R_\star$ by $\sim 5\%$ may then lead to a drop in the x-ray flux by $\lesssim 30\%$ over the descalarization timescale (~ 5 ms). Afterglow light curves in this case may appear as "broken plateaus," like that of GRB 170714A [83]. Conservation of angular momentum, however, implies that the star should spin up as a result of descalarization, and thus the drop may be less pronounced because $L_{\rm dip} \propto \Omega^4$. Likewise, the temperature of the star should increase from the compactification. Magnetohydrodynamic processes involving magnetic field reorganization may also take place, extending the dip timescale and enriching the phenomenology.

A descalarization-induced compactification may itself instigate a nuclear phase transition (e.g., quark deconfinement) due to the sudden increase in the core density [29,68]. Alternatively, the scalarized neutron star will collapse to a hairless black hole if no stable branch is available. From a scalar-GW perspective, these events would be indistinguishable [22] but could be told apart via the nature of the x-ray afterglow. If emissions persisted after the scalar energy release, a gravitational phase transition would be the favored scenario since black hole formation, which effectively terminates the stellar wind that is pumping radiation energy into the forward shock, should instead manifest as a sharp drop in the flux (as is often observed [26]).

The closest GRBs that have thus far been observed are GRBs 980425 and 170817A at distances of ~ 40 Mpc [82,84]. This distance is factor ~ 4000 times larger than that plotted in Fig. 2. As such, unless α_0 is $\gtrsim 10$ times bigger than the value we have used and year-long ($T \gtrsim$ yr) searches are carried out, we are unlikely to observe this scenario in its full capacity even with Cosmic Explorer [85] or ET [86–88] because of the $L^{-3/2}$ dependence in the effective PSD [Eq. (7)], should such stars exist. Other multimessenger possibilities for identifying a neutron star postdescalarization come from neutrino bursts (from Urca cooling or shocks triggered by compactification; cf. [89]) or indeed a burst of GWs (if the now descalarized star collapses) at some later time, either of which would again be hard to explain with a black hole remnant. It is also not necessarily the case that a neutron star must descalarize shortly after birth. Mature stars residing in the disks of active galactic nuclei or high-mass x-ray binaries [90] are particularly disposed to overaccretion. Accretion-induced collapse rates could reach $\lesssim 20\ {\rm Gpc}^{-3}\,{\rm yr}^{-1}$ from the former channel [91]. Descalarizations of Galactic stars via the latter channel should be observable with high SNR by ET. Overall however, in the absence of a detection of scalar GWs, one may not be able to tell whether a phase transition was of a nuclear or gravitational nature. This exemplifies further the well-known degeneracy between modifications of gravity and EOS uncertainty [10,34,71].

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