



Constraining Scalar-tensor Theories Using Neutron Star–Black Hole Gravitational Wave Events

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Abstract

With the continuous upgrade of detectors, greater numbers of gravitational wave (GW) events have been captured by the LIGO Scientific Collaboration and Virgo Collaboration (LVC), which offer a new avenue to test general relativity and explore the nature of gravity. Although various model-independent tests have been performed by LVC in previous works, it is still interesting to ask what constraints can be placed on specific models by current GW observations. In this work, we focus on three models of scalar-tensor theories, the Brans–Dicke theory (BD), the theory with scalarization phenomena proposed by Damour and Esposito-Farèse (DEF), and screened modified gravity (SMG). Of the four possible neutron star–black hole events that have occurred so far, we use two of them to place constraints. The other two are excluded in this work because of possible unphysical deviations. We consider the inspiral range with the cutoff frequency at the innermost stable circular orbit and add a modification of dipole radiation into the waveform template. The scalar charges of neutron stars in the dipole term are derived by solving the Tolman–Oppenheimer–Volkoff equations for different equations of state. The constraints are obtained by performing the full Bayesian inference with the help of the open source software *Bilby*. The results show that the constraints given by GWs are comparable to those given by pulsar timing experiments for DEF theory, but are not competitive with the current solar system constraints for BD and SMG theories.

Unified Astronomy Thesaurus concepts: [Gravitational waves \(678\)](#); [Non-standard theories of gravity \(1118\)](#)

1. Introduction

The theory of general relativity (GR), one of the two pillars of modern physics, is commonly regarded as the greatest theory (Chandrasekhar 1984). The splendor of GR is not only due to its elegant mathematical expression, but also its precise consistency with experimental tests. Since Einstein proposed GR in 1915, a large number of experimental tests have been conducted, ranging from laboratory scale (Adelberger 2001; Hoyle et al. 2001; Sabulsky et al. 2019) to solar system scale (Will 2014, 2018) and to cosmological scale (Jain & Khoury 2010; Clifton et al. 2012; Koyama 2016). In recent years, pulsar timing experiments (Stairs 2003; Wex 2014; Manchester 2015; Kramer 2017) and gravitational wave (GW) observations (Abbott et al. 2016a, 2019a, 2019b, 2021d) have provided great opportunities to test GR under strong field conditions. So far, all these experimental tests have supported GR at a very high level of accuracy.

Although great success has been achieved, there are still problems that GR cannot solve. At the theoretical level, GR has been facing difficulties such as singularity and quantization problems (DeWitt 1967; Kiefer 2007). At the experimental level, to explain astrophysical and cosmological observation data within the GR framework, it is necessary to introduce dark matter and dark energy whose physical nature is still unknown, which might imply the incompleteness of GR (Sahni 2004; Cline 2013). With the motivation to solve these problems, many modified gravity theories have been proposed. Among them, scalar-tensor theories are generally considered a promising candidate (Yasunori Fujii 2016).

The origin of scalar-tensor theories can be traced back to the works of Kaluza and Klein (Kaluza 1921; Klein 1926). The form that we are familiar with today was developed by works by Jordan (1955), Fierz (1956), and Brans & Dicke (1961).

Scalar-tensor theories have potential relations with dark energy, dark matter, and inflation, which continually arouse people's interest (Barrow & ichi Maeda 1990; Burd & Coley 1991; Schimd et al. 2005; Brax et al. 2006; Kainulainen & Sunhede 2006; Clifton et al. 2012). We focus on three different models of scalar-tensor theories in this work, i.e., the Brans–Dicke theory (BD), the theory with scalarization phenomena proposed by Damour and Esposito-Farèse (DEF), and screened modified gravity (SMG). The theory of Brans and Dicke (Brans & Dicke 1961) takes Mach's principle as the starting point, which says the phenomenon of inertia depends on the mass distribution of the universe. Thus the gravitational constant is promoted to be variable and coupled to the Einstein–Hilbert lagrangian as a scalar field.³ BD theory is the simplest scalar-tensor theory and is usually seen as the prototype of scalar-tensor theories, which have been well studied and constrained (Will 2018). Extensive tests have been performed in weak-field regimes based on a parameterized post-Newtonian (PPN) formalism. The most stringent constraint is given by the measurement of Shapiro time delay from the Cassini–Huygens spacecraft (Bertotti et al. 2003).

For BD theory, this tight bound requires deviations from GR in all gravitational experiments to be very small in both weak fields and strong fields. However, in the works of DEF (Damour & Esposito-Farèse 1992; Damour & Esposito-Farèse 1993), they showed that some nonperturbative effects can emerge in strong-field conditions. When the compactness of an object exceeds a critical point, spontaneous scalarization, which is usually discussed analogously with spontaneous magnetization in ferromagnets (Damour & Esposito-

³ In practice, the possible dependence of G on different circumstances is testable on a cosmological scale (Zhao et al. 2018).

Farèse 1996), arises. This phenomenon can cause the behavior in gravitational experiments involving compact objects like binary neutron star systems to have remarkable differences from experiments in weak-field regimes. In the models that can develop nonperturbative strong-field effects, order-of-unity deviations from GR are still allowed in strong-field experiments under the premise of passing the most stringent weak-field constraint. In subsequent research, different kinds of scalarization phenomena, dynamical scalarization and induced scalarization, were discovered in numerical relativity simulations of the evolution of merging binary neutron stars (Barausse et al. 2013). In binary neutron star systems, the phenomenon where the scalar field produced by the scalarized component can induce the scalarization of another component that is not scalarized initially is called induced scalarization. Since the GW event used in this work is considered a neutron star–black hole (NSBH) binary event, we do not need to be concerned with this phenomenon. Dynamical scalarization occurs in a binary system in which both components cannot be scalarized in isolation; scalarization is triggered by their gravitational binding energy of orbit. However, in previous works (Palenzuela et al. 2014; Sampson et al. 2014), it has been shown that dynamical scalarization is difficult to detect with current detectors. Therefore, we only consider spontaneous scalarization in this work.

Nonperturbative strong-field effects can be constrained by pulsar timing experiments (Damour & Esposito-Farèse 1996, 1998). Because of precise measurement technology and decades of data accumulation, the orbital period decay rate of binary pulsar systems can be measured in high precision, which makes pulsar timing experiments a good tool with which to test gravitational theories in strong-field regimes (Wex 2014). In previous works, stringent limits have been placed on gravitational theories using recent observational results from binary pulsar systems (Freire et al. 2012; Antoniadis et al. 2013; Cognard et al. 2017; Shao et al. 2017; Anderson et al. 2019; Zhao et al. 2019).

There is another class of models, SMG, which can evade tight solar system constraints by introducing screening mechanisms (Clifton et al. 2012). Various kinds of screen mechanisms have been introduced and studied, such as the Chameleon mechanism (Khouri & Weltman 2004a, 2004b), Vainshtein mechanism (Vainshtein 1972; Babichev & Deffayet 2013), and symmetron mechanism (Hinterbichler & Khouri 2010). The scalar field can be used to play the role of dark energy for driving the acceleration of the cosmic expansion on cosmological scales. Meanwhile, screening mechanisms can suppress deviations from GR on small scales to circumvent stringent constraints from the solar system tests and laboratorial experiments (see Khouri 2010, Brax 2012, Clifton et al. 2012, and Joyce et al. 2015 for comprehensive reviews). Numerous tests on SMG have also been performed in different systems (Brax et al. 2014; Liu et al. 2018a; Burrage & Sakstein 2018; Ishak 2019; Zhang et al. 2019a, 2019b, 2019c; Niu et al. 2020; Sakstein 2020).

Recently, the first GW event, GW150914, was directly detected by LIGO, which confirmed the last remaining indirectly detected prediction of GR (Abbott et al. 2016b). More GW events have been captured in subsequent observing runs by the LIGO–Virgo collaborations (LVC; Abbott et al. 2019a, 2021c). With continuing upgrades to sensitivity and the addition of new detectors, GW detections are becoming routine. GW observations offer a new avenue for testing GR

and exploring the nature of gravity in the extremely strong-field regime.

LVC has performed various model-independent tests on observed events, and no evidence for deviations from GR has been found (Abbott et al. 2016a, 2019a, 2019b, 2021d). However, for a given specific modified gravity, the model-independent parameters cannot always completely describe the deviations of GWs, which naturally depend on the characters of neutron stars and/or black holes in the corresponding theory. Therefore, it is still interesting to see what constraints on specific models can be given by current observations, which are complementary to model-independent tests. In this work, we consider the three specific scalar-tensor theories mentioned above, BD, DEF, and SMG.

Testing scalar-tensor theories using GWs has been performed since the 1990s (Will 1994). Now, greater detections of GW events and open access data allow us to constrain scalar-tensor theories using real GW data. Since in scalar-tensor gravities, the deviation of GWs from that in GR depends on the sensitivity difference of two stars, asymmetric binaries (e.g., NSBHs, white dwarf–NS binaries, white dwarf–BH binaries) are excellent targets for the model tests.

So far, among all GW events captured by LVC, there are four possible NSBH events—GW200105, GW200115, GW190426_152155, and GW190814 (Abbott et al. 2021c, 2020c, 2021a). The two events that occurred recently, GW200105 and GW200115, are the first confident observations of NSBH binaries (Abbott et al. 2021a). The component masses of these two events are consistent with current observations of black holes and neutron stars. However, the data are uninformative about the spin or tidal deformation and no electromagnetic counterparts have been detected. There is no direct evidence that the secondaries of these two events are neutron stars. Although it cannot be ruled out that the secondaries are some kind of exotic objects, we follow the most natural interpretation of these two events that they are NSBH coalescence events.

There are also two plausible NSBH events, GW190814 and GW190426_152155, in the second Gravitational Wave Transient Catalog (GWTC-2), but the nature of these two events is not definitively clear. The secondary mass of GW190814 is about $2.6M_{\odot}$, which could be interpreted as either a low-mass black hole or a heavy neutron star (Abbott et al. 2020c; Broadhurst et al. 2020; Most et al. 2020). However, according to current knowledge and observations of neutron stars, its lighter object is likely too heavy to be a neutron star (Abbott et al. 2020c). We exclude this event in our analysis. Meanwhile, the event GW190426_152155 has the highest false alarm rate (FAR; Abbott et al. 2021c). Whether it is a real signal of astrophysical origin is still not definitively clear, but its component masses are consistent with current understanding of black holes and neutron stars. There have been many recent works concerning this event, such as Li et al. (2020), Román-Garza et al. (2020), and Broekgaarden et al. (2021). Following some of them, we base our discussion on the assumption that GW190426_152155 is an NSBH coalescence event. We emphasize that our analysis is not applicable if this event is not a real NSBH binary.

There is another obstacle to our analysis. For events with a large mass ratio, deviations have been seen in the posterior distributions of the dipole modification parameter in which the GR value is excluded from 90% confidence intervals. The case of GW190814 has been shown in previous works (refer to Appendix C in Abbott et al. 2021d and Appendix A in Perkins

et al. 2021 for more detail). We have also seen similar deviations in our analysis of GW200105. The deviations are believed to be unphysical effects that are probably caused by waveform systematics, covariances between parameters, or the mode of non-GR modification parameterization. To thoroughly explain these deviations, more studies about the parameterized tests of GR on highly asymmetric sources are needed. In this work, we exclude the event GW200105 and only employ the data of GW200115 and GW190426_152155.

Zhao et al. (2019) used the binary neutron star event in GWTC-1, GW170817, to constrain scalarization effects. However, instead of directly using strain data, they employed the measurement of mass and radius from Abbott et al. (2018, 2019c) to get the constraints. In this work, we use the modification of dipole radiation in waveform and perform the full Bayesian inference to constrain scalarization effects.

The rest of this paper is organized as follows. In the next section, the modified gravity models considered in this work, including BD, DEF, and SMG, are briefly reviewed. Then, in Section 3, we present the basic information, principle of data, and statistical method used in this work. The results and conclusions are discussed in Section 4. The formulae used to get scalar charges of neutron stars by solving Tolman–Oppenheimer–Volkoff (TOV) equations are presented in Appendix A for convenience of reference. In Appendices B and C, we illustrate the comparisons of posterior distributions of other parameters with the posterior data released by LVC, and compare the constraints on the dipole radiation with the results reported by LVC. We also present the scalar charges obtained from solutions of TOV equations for all four equations of state (EoSs) considered in this work in Appendix D. A discussion of the other two possible NSBH events which are excluded in the work, GW190814 and GW200105, is presented in Appendix E. All parameter estimation samples of this work are available on Zenodo.⁴ Throughout this paper, we use the units in which $\hbar = c = 1$.

2. Scalar-tensor Theories

In this work we consider a class of scalar-tensor theories that can be described by the action

$$S = \frac{1}{16\pi G_*} \int d^4x \sqrt{-g_*} [R_* - 2g_*^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi] + S_m[\psi_m, A^2(\varphi) g_*^{\mu\nu}] \quad (1)$$

in the Einstein frame. G_* denotes the bare gravitational coupling constant, which is approximated by G when practically solving TOV equations. $g_*^{\mu\nu}$ and g_* are the Einstein-frame metric and its determinant, and $R_* \equiv g_*^{\mu\nu} R_{\mu\nu}^*$ is the Ricci scalar. The last term is the action of matter, where ψ_m collectively denotes various matter fields and $A(\varphi)$ is the conformal coupling function. Since the potential $V(\varphi)$ will be considered only in SMG theory, we do not write $V(\varphi)$ in the above action. The field equations can be derived by varying the action (1) with respect to the metric $g_*^{\mu\nu}$ and scalar field φ ,

$$R_{\mu\nu}^* = 2\partial_\mu \varphi \partial_\nu \varphi + 8\pi G_* \left(T_{\mu\nu}^* - \frac{1}{2} T^* g_{\mu\nu}^* \right), \quad \square_{g_*} \varphi = -4\pi G_* \alpha(\varphi) T_*, \quad (2)$$

where $\square_{g_*} \equiv (-g_*)^{-1/2} \partial_\mu (\sqrt{-g_*} g^{\mu\nu} \partial_\nu)$ is the curved space D'Alembertian, $T_{\mu\nu}^* \equiv 2(-g_*)^{-1/2} \delta S_m / \delta g_{\mu\nu}^*$ is the energy-momentum tensor of matter fields, and $T_* \equiv g_{\mu\nu}^* T_{\mu\nu}^*$. The quantity $\alpha(\varphi)$ is defined as $\alpha(\varphi) \equiv \partial \ln A(\varphi) / \partial \varphi$, which describes the coupling strength between the scalar field and matters. The $\ln A(\varphi)$ can be expanded around the background value φ_0 of the scalar field as

$$\ln A(\varphi) = \alpha_0(\varphi - \varphi_0) + \frac{1}{2} \beta_0(\varphi - \varphi_0)^2 + \mathcal{O}(\varphi - \varphi_0)^3, \quad (3)$$

where the coefficients α_0 and β_0 are related to two parameters β^{PPN} and γ^{PPN} in the PPN formalism by Will (2018)

$$\gamma^{\text{PPN}} - 1 = -\frac{2\alpha_0^2}{1 + \alpha_0^2}, \quad (4)$$

$$\beta^{\text{PPN}} - 1 = \frac{1}{2} \frac{\alpha_0^2 \beta_0}{(1 + \alpha_0^2)^2}. \quad (5)$$

In the context of compact binary systems, a parameter called the scalar charge, which is defined as

$$\alpha_A \equiv \left. \frac{\partial \ln m_A}{\partial \varphi} \right|_{\varphi=\varphi_0}, \quad (6)$$

can describe the coupling between the scalar field and star A. This parameter is used to determine the equation of motion and GW emission of binary systems. For compact binaries in scalar-tensor theories, the center of gravitational binding energy and the center of inertial mass are not coincident, which results in the varying dipole moment and induces extra energy loss by dipole radiation (Will 1994). We consider a GW form with the leading order of the modification, which has a dipole term in the phase (Will 1994; Zhang et al. 2017b; Liu et al. 2018b, 2020; Tahura & Yagi 2018),

$$h(f) = h_{\text{GR}}(f) \exp \left[i \frac{3}{128\eta} \varphi_{-2} (\pi G M f)^{-7/3} \right], \quad (7)$$

where φ_{-2} is given by

$$\varphi_{-2} = -\frac{5}{168} (\Delta\alpha)^2. \quad (8)$$

The constant coefficients are chosen to keep the same convention of φ_{-2} with LVC's papers (Abbott et al. 2016a, 2019a, 2019b, 2021d). $\Delta\alpha \equiv \alpha_A - \alpha_B$ is the difference between scalar charges of two bodies in a binary. For black holes, the no-hair theorem prevents them from acquiring scalar charges (Hawking 1972; Bekenstein 1995; Sotiriou & Faraoni 2012; Liu et al. 2018b). In many scalar-tensor theories, including the models considered in this work where the no-hair theorem can be applied, scalar charges of black holes are 0. For neutron stars, scalar charges can be obtained by solving TOV equations.

The detailed process of solving TOV equations to get scalar charges can be found in Damour & Esposito-Farèse (1993, 1996). We make a brief review in Appendix B for convenience of reference. Inputting the explicit form of $A(\varphi)$ and $\alpha(\varphi)$, the EoS and the initial conditions to the TOV equations, one can get the physical quantities α_A , φ_0 , and m_A

⁴ 10.5281/zenodo.4817420

outputted by Equations (A4). The coupling function $A(\varphi)$ and its logarithmic derivative $\alpha(\varphi)$ are specified by a specific theoretical model, which will be discussed in the following subsections. For the EoS, considering the constraints given by the measurement of PSR J0030+0451 (Miller et al. 2019; Riley et al. 2019) and the observational evidence that the maximum mass of a neutron star can be above $2M_\odot$ (Demorest et al. 2010; Antoniadis et al. 2013; Fonseca et al. 2016; Arzoumanian et al. 2018; Cromartie et al. 2019), we select four widely used EoSs, `sl`, `alf2`, `H4`, and `mpal`. The tabulated data of EoSs can be downloaded from the Xtreme website.⁵ To solve the differential Equations (A3), the initial conditions,

$$\mu(0) = 0, \nu(0) = 0, \varphi(0) = \varphi_c, \psi(0) = 0, \tilde{p}(0) = p_c, \quad (9)$$

need to be passed into the differential equations solver. In practice, the initial conditions are taken at the place near the center to avoid division by zero. The initial values of pressure p_c are taken on a dense grid for interpolation. The initial condition φ_c is determined by the shooting method. Different φ_c are iteratively tried until a value that can derive the desired φ_0 is found. In order to implement Monte Carlo sampling, we need to get the scalar charge at sufficient speed. It is impracticable to solve the TOV equations every time when a likelihood is evaluated. Therefore, we take the values of model parameters and p_c on a dense grid, and solve the TOV equations to get the mass and scalar charge. When the Monte Carlo sampler is running, a set of model parameters and p_c sampled by the sampler is converted to the mass and scalar charge by linear interpolation. The interpolation results will be presented in the following subsections.

2.1. Brans–Dicke Theory

We first consider the BD theory which is usually seen as the prototype of the scalar-tensor theories and has been widely studied. The BD theory is characterized by a linear coupling function given by

$$A(\varphi) = \exp(-\alpha_0 \varphi), \quad (10)$$

which led to a field-independent coupling strength $\alpha(\varphi) = \alpha_0$. There is another common convention used in the literature (Will 2014),

$$\alpha_0^2 = \frac{1}{3 + 2\omega_{\text{BD}}}. \quad (11)$$

Given the specific form of coupling function (10), we can use the process discussed in the last subsection to get the scalar charge of a neutron star. Inputting an initial condition (9) and an EoS, we can get the numerical solutions of a neutron star structure by integrating the TOV equations (A3). The quantities, α_A , φ_0 , and m_A can be extracted from the solutions by (A4). The initial condition p_c and the model parameter α_0 are taken on a dense grid for facilitating the interpolation. The last degree of freedom is the asymptotic scalar field φ_0 which is set to 0 and the initial condition φ_c is obtained by the shooting method. In order to reduce the computational burden, we use an interpolated relation $\alpha_A(\alpha_0, m_A)$ in the Monte Carlo sampling. We present the

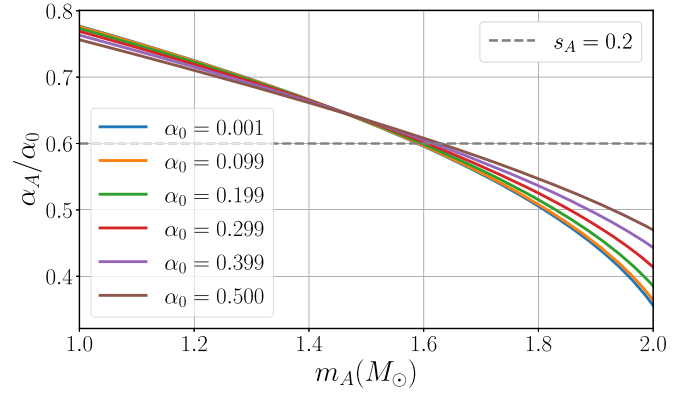


Figure 1. The interpolation result $\alpha_A(\alpha_0, m_A)$ of EoS `sl`. The gray dashed horizontal line denotes the sensitivity $s_A = 0.2$ which is an approximation commonly used in the literature.

interpolation result of EoS `sl` as an example in Figure 1 and results for other EoSs can be seen in Appendix D.

Another parameter called sensitivity, s_A , is also commonly seen in the literature. The sensitivity and the scalar charge are related by Palenzuela et al. (2014) and Sampson et al. (2014) as

$$\alpha_A = \frac{1 - 2s_A}{\sqrt{3 + 2\omega_{\text{BD}}}}. \quad (12)$$

Some works, such as Zhang et al. (2017b), employ $s_A = 0.2$ as a convenient approximation. We illustrate this approximation in Figure 1 by the gray dashed horizontal line for comparing with the results obtained by solving TOV equations.

2.2. Theory with Scalarization Phenomena

In the BD theory, all possible deviations from GR are on the order of α_0^2 in both weak-field regimes and strong-field regimes (Damour & Esposito-Farèse 1993; Will 2018). More generally, in generic scalar-tensor theories, all possible deviations from GR can be expanded as a series of powers of α_0^2 , which have the schematic form (Damour & Esposito-Farèse 1992; Esposito-Farèse 2004)

$$\text{deviation} \sim \alpha_0^2 \times \left[\lambda_0 + \lambda_1 \frac{Gm}{R} + \lambda_2 \left(\frac{Gm}{R} \right)^2 + \dots \right], \quad (13)$$

where m and R are the mass and radius of a star and $\lambda_0, \lambda_2, \dots$ are constant coefficients constructed from $\alpha_0, \beta_0, \dots$ in the expansion (3). Since solar system experiments have placed very stringent constraints on α_0 , it is plausible that all possible deviations from GR in other experiments are expected to be small. The work of DEF (Damour & Esposito-Farèse 1993) has shown that such opinions are illegitimate. In the strong-field regime, when the compactness Gm/R exceeds a critical value, some nonperturbative effects can emerge; the part in square brackets in expansion (13) can compensate for the small α_0^2 and order-of-unit deviations from GR can still be developed.

Following the model discussed by DEF (Damour & Esposito-Farèse 1993), we consider the coupling function with

⁵ http://xtreme.as.arizona.edu/NeutronStars/data/eos_tables.tar

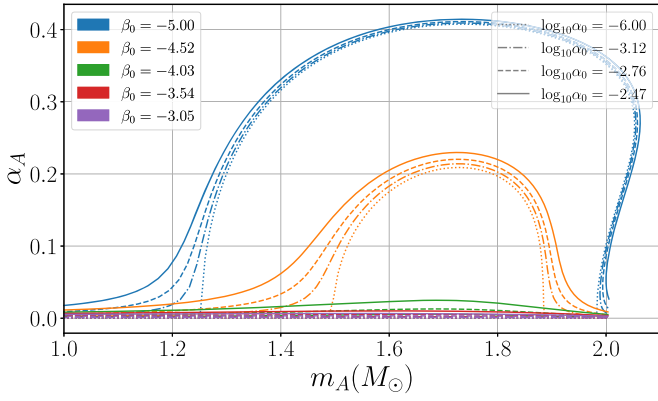


Figure 2. Nonperturbative strong-field effects in DEF. The results of one EoS, `sly`, are shown as an example. Different colors are used to denote different values of β_0 and line styles for $\log_{10} \alpha_0$. When $\beta_0 < -4$, nonperturbative strong-field effects emerge, which leads to a phase transition in a certain mass range and allows the scalar charge to be large even if the α_0 is very small. The nonzero α_0 can smooth the phase transition. The variation in the scalar charge as a function of mass is smoother for larger α_0 .

a quadratic term,

$$\ln A(\varphi) = \frac{1}{2} \beta_0 \varphi^2. \quad (14)$$

The corresponding α_0 is given by

$$\alpha_0 = -\alpha(\varphi_0) = -\beta_0 \varphi_0. \quad (15)$$

It has been shown in Damour & Esposito-Farèse (1993), when $\beta_0 < 0$, the local value of $\alpha(\varphi)$ can be amplified with respect to its asymptotic value α_0 . Nonperturbative amplification effects are expected to take place when $\beta_0 \lesssim -4$. These nonperturbative amplification effects, called spontaneous scalarization, can lead to a phase transition in a certain mass range, while, if β_0 is positive, the deviations from GR are further quenched. In this work, we only consider the negative branch. It returns to GR when $\alpha_0 = \beta_0 = 0$.

The scalar charge of a neutron star can be obtained by solving TOV equations as discussed before. We present the result of the EoS `sly` in Figure 2 as an example. There are two parameters ($\log_{10} \alpha_0, \beta_0$) characterizing the model in this case. We use colors to denote different values of β_0 , and line styles for $\log_{10} \alpha_0$. The parameter α_0 is related to weak-field effects. We consider the range of α_0 under the prior constraint of Cassini $\alpha_0 < 3.4 \times 10^{-3}$ (Bertotti et al. 2003; Damour 2009; Shao et al. 2017). Additionally, the larger α_0 can smooth the phase transition when scalarization phenomena occur (Damour & Esposito-Farèse 1996). β_0 is the parameter that controls whether spontaneous scalarization will happen in strong-field regimes. As can be seen in Figure 2, when $\beta_0 < -4$, the scalar charge can be large even if the α_0 is incredibly small for a certain mass range. The mass ranges where spontaneous scalarization can occur are different for different EoSs (Shibata et al. 2014; Shao et al. 2017). Therefore, unlike BD, the relations between the scalar charge and the mass of DEF have obvious differences for different EoSs. More details will be presented in Appendix D.

The results obtained by solving TOV equations are interpolated for stochastic sampling. Due to the scalarization phenomena, the curves representing the scalar charge as

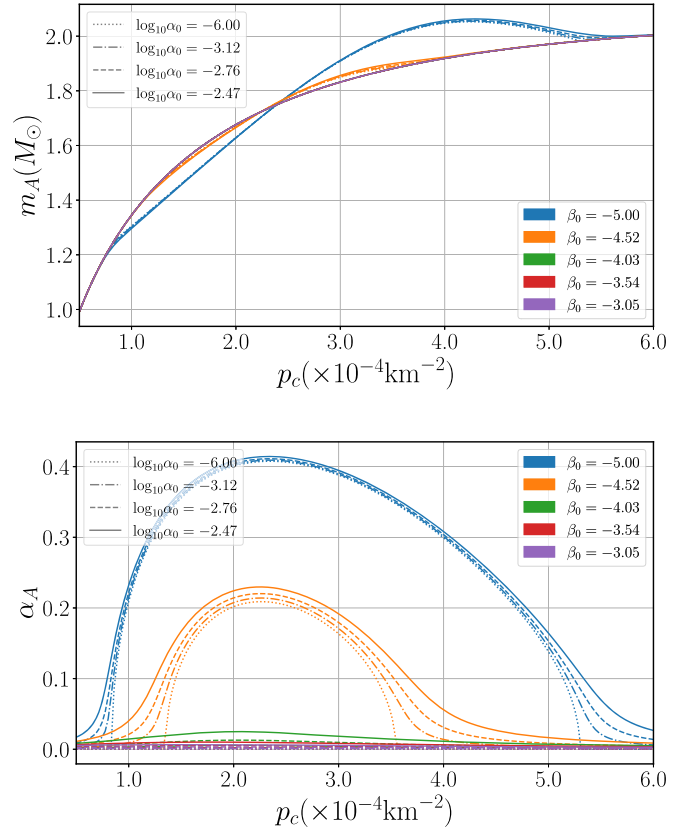


Figure 3. Interpolation results of $m_A(\log_{10} \alpha_0, \beta_0, p_c)$ and $\alpha_A(\log_{10} \alpha_0, \beta_0, p_c)$ used in stochastic sampling. We present the results of one EoS, `sly`, as an example. Since the relations between the scalar charge and the mass have hysteresis phenomena, which makes interpolation difficult, we use the parameter p_c instead of the mass in the stochastic sampling. The parameters ($\log_{10} \alpha_0, \beta_0, p_c$) are converted to parameters (α_A, m_A) which are needed to generate a waveform by the interpolation of $m_A(\log_{10} \alpha_0, \beta_0, p_c)$ and $\alpha_A(\log_{10} \alpha_0, \beta_0, p_c)$ shown in this figure.

functions of the mass have the intricate behavior of hysteresis phenomena. In order to facilitate the interpolation, instead of the mass, we use the initial condition of pressure p_c as the parameter sampled in Monte Carlo sampling and generate the interpolation function results of $\alpha_A(\log_{10} \alpha_0, \beta_0, p_c)$ and $m_A(\log_{10} \alpha_0, \beta_0, p_c)$. The results of one EoS, `sly`, are shown in Figure 3.

2.3. Screened Modified Gravity

The third model we considered is SMG. Besides the coupling function $A(\varphi)$ characterizing the interaction between the scalar field and the matter field, there is the potential $V(\varphi)$ characterizing the self-interaction of the scalar field. The coupling function $A(\varphi)$ and the potential $V(\varphi)$ define the effective potential $V_{\text{eff}}(\varphi)$ which controls the behavior of the scalar field. The scalar field acquires the mass around the minimum of the effective potential $V_{\text{eff}}(\varphi)$, which depends on the environmental density. The mass of the scalar field can be large in high-density regions and the range of the fifth force becomes short, so the effects of the scalar field are screened. While, on large scales, the environmental density is low, the scalar field becomes light and can affect the galactic dynamic or the universe expansion acceleration (see the comprehensive review by Ishak 2019 for additional types of screening mechanisms). For the general SMG with canonical kinetic

Table 1
Basic Information on Six Events Likely Including at Least One Neutron Star

Event	Type	$m_1(M_\odot)$	$m_2(M_\odot)$	SNR	FAR (yr ⁻¹)
GW170817	BNS	$1.46^{+0.12}_{-0.10}$	$1.27^{+0.09}_{-0.09}$	33.0	$\leq 1.0 \times 10^{-7}$
GW190425	BNS(?)	$2.0^{+0.6}_{-0.3}$	$1.4^{+0.3}_{-0.3}$	13.0	7.5×10^{-4}
GW190426_152155*	NSBH(?)	$5.7^{+3.9}_{-2.3}$	$1.5^{+0.8}_{-0.5}$	10.1	1.44
GW190814	NSBH(?)	$23.2^{+1.1}_{-1.0}$	$2.6^{+0.08}_{-0.09}$	22.2	$\leq 1.0 \times 10^{-5}$
GW200105	NSBH	$8.9^{+1.2}_{-1.5}$	$1.9^{+0.3}_{-0.2}$	13.9	0.36
GW200115*	NSBH	$5.7^{+1.8}_{-2.1}$	$1.5^{+0.7}_{-0.3}$	11.6	$\leq 1.0 \times 10^{-5}$

Note. The data are copied from the Gravitational Wave Open Science Center (<http://www.gw-openscience.org>). The two events with stars are used to place the constraints in this work.

energy term, we can rewrite the action as

$$S = \int d^4x \sqrt{-g_*} \left[\frac{1}{16\pi G} R_* - \frac{1}{2} g_*^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi - V(\varphi) \right] + S_m[\psi_m, A^2(\varphi) g_*^{\mu\nu}], \quad (16)$$

where bare potential $V(\varphi)$ characterizes the scalar self-interaction, which endows the scalar field with mass. There are many SMG models in the market, including the chameleon, symmetron, dilaton, and $f(R)$ models, in which the functions $V(\varphi)$ and $A(\varphi)$ are chosen as the specific forms (Zhang et al. 2016; Liu et al. 2018a). The scalar field equation of motion can be yielded by varying the action with respect to φ ,

$$\square_{g_*} \varphi = \frac{\partial}{\partial \varphi} V_{\text{eff}}(\varphi), \quad (17)$$

where the effective potential is defined as

$$V_{\text{eff}}(\varphi) = V(\varphi) - T_*. \quad (18)$$

The waveform of GWs from inspiraling compact binaries in the SMG has been given in the previous work (Liu et al. 2018b). As mentioned above, we only consider the leading order modification which is the dipole term shown in Equations (7) and (8). Since the effects of the scalar field are suppressed due to the screening mechanism, the scalar charges of neutron stars are expected to be small. Therefore, we do not solve the TOV equations to get the scalar charge, but adopt a simple approximation that considers a neutron star as a static spherical symmetric object with constant density. The scalar field Equation (17) can be simplified and solved directly to get the exact solution. Matching the internal and external solutions, the scalar charge of a neutron star in the SMG can be given by (see Appendix A in Zhang et al. 2017a for more details)

$$\alpha_A = \frac{\varphi_{\text{VEV}}}{M_{\text{Pl}} \Phi_A}, \quad (19)$$

where $M_{\text{Pl}} = \sqrt{1/8\pi G}$ is the reduced Planck mass, φ_{VEV} is the vacuum expectation value of the scalar field, and $\Phi_A = Gm/R$ is the surface gravitational potential of object A.

3. Public Data and Bayesian Method

3.1. Public GW Data

Among all GW events released by LVC, there are four possible NSBH events (Abbott et al. 2021c, 2020c, 2021a), GW190426_152155, GW190814, GW200105, GW200115, and two (possible) binary neutron star (BNS) events

(Abbott et al. 2017a, 2020a), GW170817, GW190425, which could potentially be used for tests of scalar-tensor theories.

For convenience of reference, we list some basic information of these six events in Table 1 and present a brief review of these events below. GW170817 (Abbott et al. 2017a) is a relatively confident BNS event since its electromagnetic counterpart was captured by various facilities across the electromagnetic spectrum (Abbott et al. 2017b). While definite electromagnetic counterpart observations for the all other events are absent. For GW190425 (Abbott et al. 2020a), the mass of its components is consistent with neutron stars, but its total mass and chirp mass are larger than those of any other known binary neutron star systems. It cannot be ruled out by GW data alone that one or both of its components are black holes. GW190814 (Abbott et al. 2020c) is a stranger event with its significantly unequal mass ratio and unusual secondary component. It involves a $22.2\text{--}24.3M_\odot$ black hole and a $2.50\text{--}2.67M_\odot$ object which we do not know much about yet. All current models of formation and mass distribution for compact binaries are challenged by this event. GW190426_152155 (Abbott et al. 2021c) is a possible NSBH event since the mass of its components is consistent with our current understanding of neutron stars and black holes. But this event has the highest FAR, 1.4 yr^{-1} , which casts doubt on whether it is a real signal of astrophysical origin. Additionally, since the data are uninformative about the effects such as tidal deformability or spin-induced quadrupole, it is not possible to rule out that its secondary object is a black hole or other exotic object. GW200105 and GW200115 are two NSBH coalescence events reported recently (Abbott et al. 2021a). The primaries and secondaries of these two binaries have masses within the range of known black holes and neutron stars, respectively. These two events have been regarded as the first observations of NSBH binaries via any observational means. Note that, although the most natural interpretations of these two events are NSBH coalescences, this conclusion is inferred only by their component masses. Until now, there is no direct evidence, such as tidal or spin deformation and electromagnetic counterparts. It is still difficult to rule out whether the secondaries are other objects.

Although there are six events that probably include at least one neutron star in the current GW catalog, only two events, GW190426_152155 and GW200115, can be used in this work.

For GW190814, due to its unusual mass ratio which is in a region that has not been systematically studied, the issue of waveform systematics can lead to some kind of unphysical deviation (refer to Appendix C in Abbott et al. 2021c and Appendix A in Perkins et al. 2021 for more details). To our knowledge, there are no EoSs that can reach the mass of its

secondary object and also be favored by current observations of neutron stars. Thus, we exclude this event in our discussion.

One of two NSBH events reported recently, GW200105, also has a large mass ratio, which can be seen in Figure 14. A similar unphysical deviation is also present in the analysis of this event. We show the posterior distribution of the dipole modification parameter in Figure 13. The GR value falls in the tail of the posterior and is excluded from the 90% confidence interval. This deviation is believed to be unphysical, which might be the consequence of systematic errors of waveform templates, covariances between parameters, or the parameterization of non-GR modification (Abbott et al. 2021c; Perkins et al. 2021). We present further discussion on this issue in Appendix E and exclude this event in the main body of this work.

For GW170817 and GW190425, only one side limit can be placed on the mass ratio, which means a situation where the two components have an equal mass cannot be ruled out. The dipole radiation depends on the difference of the scalar charges between two components of a binary. As shown in Figures 1 and 2, the scalar charges are functions of mass for BD $\alpha_A(\alpha_0, m)$ and DEF $\alpha_A(\alpha_0, \beta_0, m)$. Symmetrical binaries can lead to very long tails in posterior distributions of α_0 or (α_0, β_0) which cannot descend to zero when reaching the boundary of a prior setting. Therefore, even though the two BNS events can place very strong bounds on the dipole amplitude, we cannot use them to place any effective constraints on model parameters of BD or DEF. But we can constrain the dipole radiation for GW170817 and GW190425 without considering specific model parameters. In order to compare with the results from LVC, we also perform the constraints on φ_{-2} for these two events in Appendix C.

Although the origin of GW190426_152155 still has some uncertainty, the data are consistent with a GW signal from an NSBH coalescence. We think it is feasible to test modified gravity models using this event. The results can at least offer a reference for future, more confident NSBH events.

As discussed above, the events GW190426_152155 and GW200115 are the only two left that can be used for our purpose. The data are downloaded from the GW Open Science Center⁶ (Abbott et al. 2021b) and downsampled to 2048 Hz. Besides strain data, power spectral densities (PSDs) are also needed for parameter estimation (Abbott et al. 2020b). Instead of directly estimating PSDs from strain data by the Welch method, we use the event-specific PSDs, which are encapsulated in LVC posterior sample releases for specific events (LIGO Scientific Collaboration & Virgo Collaboration 2020a, 2020c). These PSDs are expected to lead to more stable and reliable parameter estimation (Cornish & Littenberg 2015; Littenberg & Cornish 2015; Abbott et al. 2019a). As mentioned above, we only consider inspiral stages; therefore, the frequency corresponding to the innermost stable circular orbit (ISCO),

$$f_{\text{ISCO}} = \frac{1}{6^{3/2}\pi M}, \quad (20)$$

where M denotes the total mass of the binary, is chosen as the maximum frequency cutoff (Buonanno et al. 2009). The minimum frequency cutoffs are chosen by following LVC's papers (Abbott et al. 2021c, 2021a). The frequency of GW from inspiral of a compact binary in circular orbit evolves with time. The data segment durations are set to be consistent with

this frequency range. The data segment is positioned such that there is a 2 s post-trigger duration (Romero-Shaw et al. 2020).

3.2. Bayesian Method

Bayesian inference is broadly used in modern science for extracting useful information from noisy data. Bayesian inference allows us to make statements on how probabilities of parameters distribute in a priori ranges based on the observed data in a specific model. In the context of GW astronomy, given a model M described by a set of parameters θ , observed strain data d , and background information I which determines the likelihood and prior, the Bayes' theorem can be written as (Bayes 1763; Abbott et al. 2020b)

$$p(\theta|d, M, I) = p(\theta|M, I) \frac{p(d|\theta, M, I)}{p(d|M, I)}. \quad (21)$$

The left-hand side is the posterior probability density function of model parameters, which is the product of Bayesian inference and represent the result inferred from data. The three terms on the right-hand denote the prior probability density $p(\theta|M, I)$, the likelihood $p(d|\theta, M, I)$, and the evidence $p(d|M, I)$. Under the assumption that the noise from detectors is stationary and Gaussian, the likelihood function can be written as (Cutler & Flanagan 1994; Romano & Cornish 2017)

$$p(d|\theta, M, I) \propto \exp \left[-\frac{1}{2} \sum_i \langle h(\theta) - d | h(\theta) - d \rangle \right], \quad (22)$$

where i denotes different detectors and $h(\theta)$ is the waveform template. The angle brackets represent the noise-weighted inner product defined as

$$\langle a|b \rangle = 4\Re \int \frac{a(f)b^*(f)}{S_n(f)} df \quad (23)$$

with the noise PSD $S_n(f)$ of the detector.

As discussed in Section 2, we consider a waveform model including a term of dipole radiation. The waveform template used to compute likelihood is obtained by slightly modifying the aligned spin with a tidal deformability waveform IMRPhenomD_NRTidal (Dietrich et al. 2019) which has been implemented in the LIGO Algorithm Library LALSuite (LIGO Scientific Collaboration 2018).

For the prior, a range needs to be set for each parameter of the model. As discussed above, instead of the mass parameter, we choose the central pressure of a neutron star as a model parameter. The prior ranges are set by referring to Romero-Shaw et al. (2020) and Abbott et al. (2019a). According to the known properties of binary neutron stars, we employ the low spin prior in this work (Burgay et al. 2003; Stovall et al. 2018; Abbott et al. 2019c). The evidence plays the role of normalization factor and is also used in model selection.

One of the obstacles to applying Bayesian inference is the extremely costly computation. For the huge parameter space, it is impractical to evaluate the likelihood on a grid. The Markov Chain Monte Carlo (MCMC) methods (Metropolis et al. 1953; Hastings 1970) or nested sampling methods (Skilling 2004, 2006) are commonly used to estimate the

⁶ <https://doi.org/10.7935/99gf-ax93>

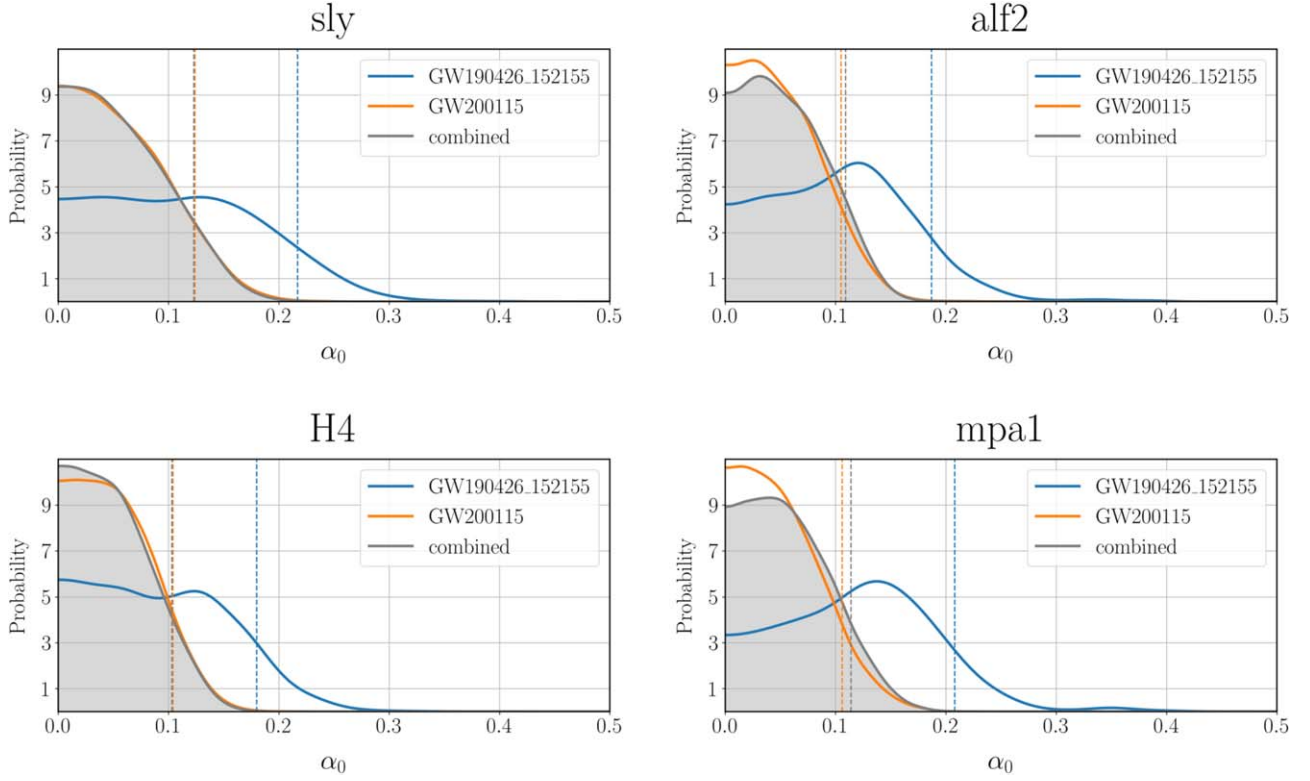


Figure 4. Posterior distributions of α_0 for BD. The results of two events are shown by blue and orange lines. The gray lines with translucent shading denote the combined posterior distributions. The dashed vertical lines indicate the upper limits at 90% CL.

posterior distribution by sampling in parameter space. We use the open source library *Bilby*⁷ (Ashton et al. 2019) with the nested sampler *Dynesty*⁸ (Speagle 2020) to do our Bayesian inference. The sampler settings are chosen by referring to Romero-Shaw et al. (2020).

4. Results and Conclusions

We will present our results in this section. All our results are consistent with GR. For the parameter β_0 in DEF, we find the constraints given by GWs are comparable with the previous constraints given by pulsar timing experiments. For BD and SMG, the constraints are not competitive with the current bounds placed by solar system experiments. We do not find significant differences among the constraints using different EoSs. More details follow.

4.1. Brans–Dicke Theory

For BD, the posterior distributions of α_0 are shown in Figure 4. The posteriors of two events can be combined (Agathos et al. 2014; Abbott et al. 2019d), and the combined results are shown by the gray lines with translucent shading. The vertical dashed lines denote the upper limits of α_0 at 90% confidence level (CL) whose exact values are collected in Table 2. Colors are used to denote two events. In the results, the impact of different EoSs is invisible within statistical errors. According to the relation (11), one can get the constraints on ω_{BD} which are also shown in Table 2. So far, the most stringent constraint on BD is from the measurement of Shapiro time delay performed by the Cassini spacecraft which places the

Table 2

90% CL Limits of Combined Posteriors for the Model Parameter, α_0 , in BD and its Corresponding ω_{BD} , as Well as the Parameter β_0 in DEF

	BD		DEF
	α_0	ω_{BD}	β_0
sly	$\gtrsim 0.123$	$\gtrsim 31.5$	$\gtrsim -3.93$
alf2	$\gtrsim 0.109$	$\gtrsim 40.6$	$\gtrsim -4.00$
H4	$\gtrsim 0.103$	$\gtrsim 45.6$	$\gtrsim -3.77$
mpa1	$\gtrsim 0.114$	$\gtrsim 37.0$	$\gtrsim -4.08$

bound (Bertotti et al. 2003),

$$\gamma^{\text{PPN}} - 1 = (2.1 \pm 2.3) \times 10^{-5}. \quad (24)$$

The corresponding constraint on ω_{BD} is (Will 2014)

$$\omega_{\text{BD}} > 40000. \quad (25)$$

The pulsar timing experiments also place the constraint (Freire et al. 2012; Antoniadis et al. 2013; Zhang et al. 2019a)

$$\omega_{\text{BD}} > 13,000. \quad (26)$$

We summarize the different constraints in Table 3 for comparison. The constraints given by GWs cannot compete with the current constraints. This result was expected. In Zhang et al. (2017b), we found that in the third-generation GW detector era, the bound by combining a larger number of GW events is expected to be better than that derived in solar system.

4.2. Theory with Scalarization Phenomena

For DEF, we plot the posterior distributions of $(\log_{10} \alpha_0, \beta_0)$ in Figure 5 and summarize the combined constraints of parameter β_0 in Table 2. In Figure 5, we show the 90% CL regions of the joint

⁷ <https://github.com/lscsoft/bilby>

⁸ <https://github.com/joshspiegel/dynesty>

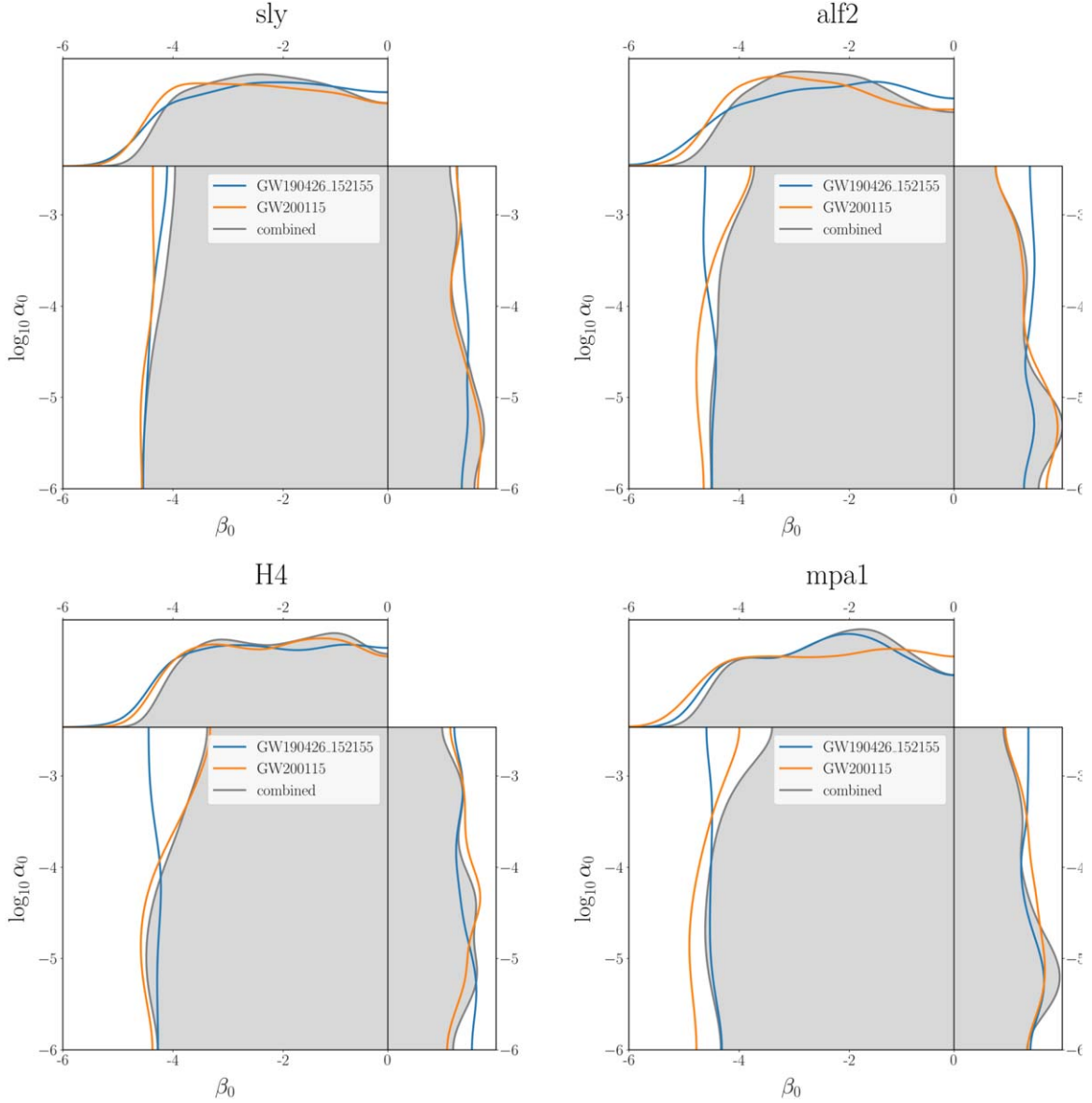


Figure 5. The posterior distributions of parameters ($\log_{10} \alpha_0$, β_0) in DEF. The 90% CL regions of the joint posterior distributions for ($\log_{10} \alpha_0$, β_0) are shown in the main panels and the marginalized posteriors for $\log_{10} \alpha_0$ and β_0 are plotted in the side panels. The posterior distributions of two events are indicated by two different colors and the combined posteriors are shown by the gray lines with translucent shading. GR returns when $\alpha_0 = \beta_0 = 0$. The results show consistency with GR and no evidence for scalarization phenomena. In the range of $\beta > -4$, scalarization cannot occur and the scalar charges are almost zero. The different values of β_0 can hardly be distinguished. Therefore, the distributions of β_0 are flat in this range. Since we require the prior of $\log_{10} \alpha_0$ to be compatible with the Cassini constraint and the influence on the scalar charges of $\log_{10} \alpha_0$ is too small to be detected, the different values of $\log_{10} \alpha_0$ in its prior range are totally indistinguishable by the stochastic sampler.

Table 3
Different Constraints Summarized for Comparison

	Solar System	Pulsar Timing	GWs (combined)	GWs (only GW200115)
BD	$\omega_{\text{BD}} \gtrsim 40000$	$\omega_{\text{BD}} \gtrsim 13000$	$\omega_{\text{BD}} \gtrsim 40$	$\omega_{\text{BD}} \gtrsim 40$
DEF	$\dots$	$\beta_0 \gtrsim -4.3$	$\beta_0 \gtrsim -4.0$	$\beta_0 \gtrsim -4.2$
SMG	$\frac{\varphi_{\text{VEV}}}{M_{\text{Pl}}} \lesssim 7.8 \times 10^{-15}$	$\frac{\varphi_{\text{VEV}}}{M_{\text{Pl}}} \lesssim 4.4 \times 10^{-8}$	$\frac{\varphi_{\text{VEV}}}{M_{\text{Pl}}} \lesssim 1.8 \times 10^{-2}$	$\frac{\varphi_{\text{VEV}}}{M_{\text{Pl}}} \lesssim 1.8 \times 10^{-2}$

Note. We also list the results from event GW200115 only, in case the event GW190426_152155 is believed to be a false GW signal. However, for BD and SMG, differences between the combined results and the results excluding GW190426_152155 are within the round-off errors.

posterior distributions for $(\log_{10} \alpha_0, \beta_0)$ in the main panels and the marginalized posteriors for $\log_{10} \alpha_0$ and β_0 are plotted in the side panels. The blue and orange lines denote the two events, respectively, and gray lines with translucent shading denote the combined results.

Although the mass ranges where the scalarization can occur are different for different EoSs (Shibata et al. 2014; Shao et al. 2017), we do not find the results have obvious differences beyond statistical errors for different EoSs. GR returns when $\alpha_0 = \beta_0 = 0$. Our results are consistent with GR and have no evidence for scalarization phenomena. There are some features in Figure 5 which might be noteworthy.

The posterior distributions of β_0 are almost flat when $\beta_0 > -4$. This is because the scalarization phenomena cannot occur in this range. As can be seen in Figure 2 and the bottom panel of Figure 3, in the range of $\beta_0 > -4$, the nonperturbative effects will not take place for any neutron star mass. The scalar charges are almost independent of β_0 . Different values of β_0 can hardly be distinguished by the sampling algorithm. Hence, the posterior distributions are flat in this range.

The posteriors of $\log_{10} \alpha_0$ distribute uniformly on the prior range, which shows no difference with the prior distribution. The two-dimensional joint distribution in the main panel of Figure 2 also shows that different values of $\log_{10} \alpha_0$ are totally indistinguishable for the stochastic sampler. On the one hand, in the range of $\beta_0 > -4$ where nonperturbative amplification effects cannot occur, it resembles the case of BD theory. We adopt a prior range compatible with the Cassini constraint, in which the values of α_0 are incredibly small. Any scalar charges evaluated in this region are too small to cause detectable effects. Different values of β_0 and $\log_{10} \alpha_0$ cannot be distinguished by the sampler in this region. On the other hand, even in the range of $\beta_0 < -4$, as can be seen in Figures 2 and 3, the influence of $\log_{10} \alpha_0$ on the scalar charges is much smaller than β_0 . The small difference caused by $\log_{10} \alpha_0$ cannot be detected by the noisy GW data. For these reasons, we cannot place constraints on the parameter $\log_{10} \alpha_0$ from our sampling results.

It is useful to compare our results with those of previous similar works (Shao et al. 2017; Zhao et al. 2019) which used pulsar timing experiments to constrain DEF. In Zhao et al. (2019), GW event GW170817 is also considered to place constraints. Different from the full Bayesian method in which the waveform templates and the PSD are used to construct the likelihood function, they employed the measurement results of mass and radii to construct the likelihood. Another difference is that we use the prior range of $-6 < \beta_0 < 0$ which can return to GR at the edge. Shao et al. (2017) and Zhao et al. (2019), were only interested in the range $\beta_0 \in [-5, -4]$ where scalarization can take place. They present the 90% CL bounds

$$\beta_0 \gtrsim -4.3. \quad (27)$$

Our constraints of β_0 are better on the order of 0.3. However, considering the statistical errors, we think this difference is not significant. The different prior setting could also induce this slight difference. For α_0 , this could place the constraint $\alpha_0 \lesssim 10^{-4}$, while the different values of $\log_{10} \alpha_0$ are indistinguishable in our sampling. As discussed above, due to the statistical uncertainty and the reason that we consider the prior of β_0 including the range where scalarization cannot occur, we cannot constrain $\log_{10} \alpha_0$. As can be seen in Figure 15 of Appendix A in Zhao et al. (2019), the parameter $\log_{10} \alpha_0$ also cannot be constrained well by using the GW only. The

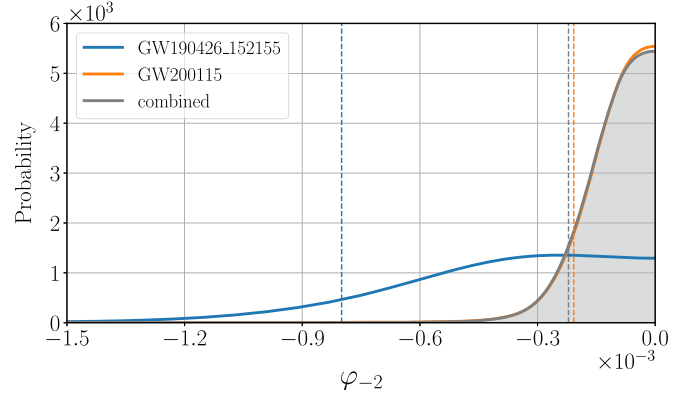


Figure 6. The posterior distribution of φ_{-2} . As mentioned in the main text, we only consider the physically possible prior range, $\varphi_{-2} \leq 0$. The blue and orange colors are used to denote the two events and the combined result is indicated by the gray color. The dashed vertical lines denote the limits at 90% CL.

constraints given by GW170817 in Zhao et al. (2019) are a little more related to EoS compared to our results. This is because of the different mass parameters of two GW events. The primary and secondary mass of GW170817 with the low-spin prior assumption at 90% CL are given by Abbott et al. (2018, 2019c) as $m_1 \in (1.36, 1.60)$ and $m_2 \in (1.16, 1.36)$. As can be seen in Figure 12, the scalarization phenomena on these ranges depend more strongly on the EoS, while the secondary masses of the two events considered here are heavier and in the range where the dependence of scalarization phenomena on EoS is less. We summarize the comparisons in Table 3.

4.3. Screened Modified Gravity

The third model we discuss is SMG. As mentioned in Section 2, the screening mechanism can suppress the effects of the scalar field in high-density regions. The scalar charges of neutron stars are expected to be small. Hence, for SMG we do not consider different EoS and strictly solve the TOV equations but adopt a simple approximation that considers neutron stars have a constant density in order to get the scalar charges as presented in Equation (19). We use the typical values $m = 1.4M_\odot$ and $R = 10$ km for the surface gravitational potential Φ_A in Equation (19). Since this scalar charge is independent with other parameters under the approximation, we do not sample parameters of specific SMG models whereas we sample the parameter φ_{-2} in Equation (7) and place the constraint on φ_{VEV} by the upper limit of φ_{-2} . Constraining the parameter φ_{-2} is similar to the model-independent parameterized tests of GW generation performed by LVC (Abbott et al. 2019a, 2019b, 2021d) except for two differences. Since we are discussing the specific model, it is more logical to take physical limits into consideration. For the models considered in this works, the dipole radiation always takes energy away and the outgoing energy flux is positive. The phase evolution will be ahead compared with the case of GR. So, we consider the prior range constraining $\varphi_{-2} \leq 0$. Another difference is that we only consider the inspiral range. Since we are ignorant about the waveform in the merge and ringdown range for scalar-tensor theories, we adopt the cutoff at the frequency corresponding to ISCO as shown in the Equation (20).

The posterior distribution of φ_{-2} is shown in Figure 6. The combined constraint at 90% CL is

$$\varphi_{-2} > -2.2 \times 10^{-4}, \quad (28)$$

and the corresponding constraint on φ_{VEV} is given by

$$\frac{\varphi_{\text{VEV}}}{M_{\text{Pl}}} < 1.8 \times 10^{-2}. \quad (29)$$

The constraint on φ_{-2} by GW170817 is about 10^{-5} (Abbott et al. 2019b) which is one order of magnitude better than the constraint given here. This better constraint is because there are more circles that can be monitored for GW170817. Since GW170817 is lighter than the two events considered here, in the detectors' sensitive band, the signal can be observed for longer and the circles can be tracked more. We show the posterior distributions of φ_{-2} given by GW170817 and another possible binary neutron star event GW190425 in Appendix C for convenience of comparison.

The parameter φ_{VEV} also has to be constrained by the solar system tests and pulsar timing experiments. The most stringent constraint in the solar system is from the lunar laser ranging (LLR) measurement (Hofmann et al. 2010; Zhang et al. 2019a), which is given by

$$\frac{\varphi_{\text{VEV}}}{M_{\text{Pl}}} < 7.8 \times 10^{-15}. \quad (30)$$

Pulsar timing experiments also place the constraint (Freire et al. 2012; Antoniadis et al. 2013; Zhang et al. 2019a)

$$\frac{\varphi_{\text{VEV}}}{M_{\text{Pl}}} < 4.4 \times 10^{-8}. \quad (31)$$

These constraints are much better than the constraint obtained in this work. On one hand, these much stronger constraints are caused by the fact that the surface gravitational potentials of white dwarfs and objects in the solar system are much less than those in neutron stars. The difference of compactness between white dwarfs and neutron stars can be about $\Phi_{\text{WD}}/\Phi_{\text{NS}} \sim 10^{-4}$. This ratio will be much less for objects in the solar system. On the other hand, after the GW signal enters the sensitive band, there are only tens of seconds left before the final plunge, whereas the pulsar timing experiments can monitor the orbital motion of a binary at lower frequency and for a longer time. Experiments in the solar system can also collect data over long periods of time.

5. Summary

As more and more various kinds of GW events are observed, GW is becoming an important tool to test GR and explore the nature of gravity. The open access data and user-friendly software tools engage the community to take part in the research about GWs more broadly. Although various model-independent tests have been performed by LVC and have placed stringent upper limits on possible deviations from GR, it is still interesting to ask what constraints can be placed on specific models by recent observations. In this work, we consider three specific scalar-tensor theories, the BD theory, the theory with scalarization phenomena proposed by DEF, and the SMG.

The data used in this work are the possible NSBH coalescence GW event GW190426_152155 in GWTC-2 and one of the two confident NSBH events reported recently, GW200115. Due to possible unphysical deviations, we exclude

the events GW190814 and GW200105 in this work. Since the dipole amplitude depends on the difference between the scalar charges of two components of a binary, if the possibility that the two components have an equal mass cannot be ruled out, we are unable to place an effective constraint. Therefore, we also exclude the two BNS events GW170817 and GW190425 in the analysis.

We place constraints by performing the full Bayesian inference. The waveform template with the dipole term which is the leading order of modification is used to construct the likelihood. The dipole radiation in scalar-tensor theories is proportional to the square of the scalar charge difference between two component objects of a binary. The scalar charges of black holes are zero which is assured by the no-hair theorem. The scalar charges of neutron stars are obtained by solving TOV equations for BD and DEF. For SMG, the effects of the scalar field are expected to be small due to the screening mechanism. So, we adopt a simple assumption that the density of a neutron star is a constant to get the scalar charge.

Four tabular EoSs are used when solving TOV equations to get the scalar charges for BD and DEF. However, we do not find the different EoS have significant influence on the results. All results we get are consistent with GR. The constraint on BD is about $\alpha_0 \lesssim 0.1$ or equivalent to $\omega_{\text{BD}} \gtrsim 40$. For DEF, we get the constraint $\beta_0 \gtrsim -4.0$. Due to our prior settings and statistical uncertainties, we cannot get the constraints of the parameter $\log_{10} \alpha_0$ in DEF. For SMG, we place the upper limit $\varphi_{\text{VEV}}/M_{\text{Pl}} \lesssim 1.8 \times 10^{-2}$. All constraints presented above are at 90% confidence. The constraint on β_0 in DEF is comparable with the previous constraint from pulsar timing experiments. The constraints on BD and SMG have no comparison with previous constraints given by the solar system tests and pulsar timing experiments. Although the results of this work do not find any new phenomena or push the current constraints to be more stringent, our results complement the tests on these three specific models in the strong-field regime and make preparations for future NSBH events.

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Facilities: LIGO, Virgo.

Software: Bilby (Ashton et al. 2019), Dynesty (Speagle 2020), LALSuite (LIGO Scientific Collaboration 2018), PESummary (Hoy & Raymond 2021), NumPy (van der Walt et al. 2011; Harris et al. 2020), SciPy (Virtanen & Gommers et al. 2020), matplotlib (Hunter 2007).

Appendix A

Differential Equations for Neutron Star Structure

The scalar charge of a neutron star can be obtained by solving the TOV equations. The TOV equations for a neutron star in scalar-tensor theories can be found in previous works (Damour & Esposito-Farèse 1993, 1996). We present a succinct summary here for the convenience of reference. Assuming that the neutron star is isolated and nonrotating, the geometry part can be given by the static spherically symmetric metric

$$ds_*^2 = g_{\mu\nu}^* dx^\mu dx^\nu = -e^{\nu(r)} dt^2 + \frac{dr^2}{1 - 2\mu(r)/r} + r^2(d\theta^2 + \sin^2\theta d\varphi^2). \quad (\text{A1})$$

The matter part is described by the perfect-fluid form of energy-momentum tensor in the Jordan frame

$$\tilde{T}^{\mu\nu} = (\tilde{\rho} + \tilde{p})\tilde{u}^\mu\tilde{u}^\nu + \tilde{p}\tilde{g}^{\mu\nu}. \quad (\text{A2})$$

We use tilde to denote a quantity in the Jordan frame and star to denote a quantity in the Einstein frame. $\tilde{T}$ and T_* are related by $T_* = A^4(\varphi)\tilde{T}$. Taking the above metric (A1) and energy-momentum tensor (A2) into the field Equations (2) and energy-momentum conservation equation $\tilde{\nabla}_\mu \tilde{T}^{\mu\nu} = 0$, one can get the following differential equations, which describe the structure of

neutron star,

$$\begin{aligned} \mu' &= 4\pi G_* r^2 A^4(\varphi) \tilde{\rho} + \frac{1}{2} r(r - 2\mu) \psi^2 \\ \nu' &= 8\pi G_* A^4(\varphi) \frac{r^2}{r - 2\mu} \tilde{p} + r\psi^2 + \frac{2\mu}{r(r - 2\mu)} \\ \varphi' &= \psi \\ \psi' &= 4\pi G_* A^4(\varphi) \frac{r}{r - 2\mu} [\alpha(\varphi)(\tilde{\rho} - 3\tilde{p}) + r\psi(\tilde{\rho} - \tilde{p})] \\ &\quad - \frac{2(r - \mu)}{r(r - 2\mu)} \psi \\ \tilde{p}' &= -(\tilde{\rho} - \tilde{p}) \left[4\pi G_* \frac{r^2 A^4(\varphi) \tilde{p}}{r - 2\mu} + \frac{1}{2} r\psi^2 + \frac{\mu}{r(r - 2\mu)} \right. \\ &\quad \left. + \alpha(\varphi)\psi \right]. \end{aligned} \quad (\text{A3})$$

The above equations can be solved once the EoS, which is the relation between $\tilde{\rho}$ and $\tilde{p}$, and the initial conditions are given. Physical quantities, the scalar charge, the scalar field at infinity, and the gravitational mass, can be extracted from the solution by matching the interior and exterior solutions,

$$\alpha_A = -\frac{2\psi_s}{\nu'_s} \quad (\text{A4})$$

$$\varphi_0 = \varphi_s + \frac{2\psi_s}{(\nu'_s{}^2 + 4\psi_s^2)^{1/2}} \operatorname{arctanh} \left[\frac{(\nu'_s{}^2 + 4\psi_s^2)^{1/2}}{\nu'_s + 2/r_s} \right] \quad (\text{A5})$$

$$\begin{aligned} m_A &= \frac{r_s^2 \nu'_s}{2G_*} \left(1 - \frac{2\mu_s}{r_s} \right)^{1/2} \\ &\quad \times \exp \left[-\frac{\nu'_s}{(\nu'_s{}^2 + 4\psi_s^2)^{1/2}} \operatorname{arctanh} \left(\frac{(\nu'_s{}^2 + 4\psi_s^2)^{1/2}}{\nu'_s + 2/r_s} \right) \right] \end{aligned} \quad (\text{A6})$$

where the subscript s denotes that the quantities take the values at the star surface.

Appendix B

Posterior Distribution of Other Parameters

In order to verify the reliability of our sampling, we compare our results with the posterior data released by LVC.⁹ We select one of our multiple runs for each event as an example to plot together with parameter estimation samples in the posterior data files released by LVC in Figures 7 and 8. The posterior distributions of some interior parameters and the luminosity distance are presented by the corner plot.

The definitions and labels of the parameters follow the conventions implemented in `bilby`. The blue lines and

regions denote our results and the red for the results from LVC. The dashed vertical lines represent the 5% and 95% quantiles. Since we make our discussion based on the assumption that the secondary of GW190426_152155 is a neutron star and impose a constraint $m_2 \in [1.0, 2.0]$, the result of mass ratio has slight differences with the result from LVC. Due to the degeneracy between aligned spin and mass ratio, the result of effective inspiral spin parameter has also a little mismatch with LVC result. Except for this, the sampling results of other parameters are consistent with the results released by LVC quite well.

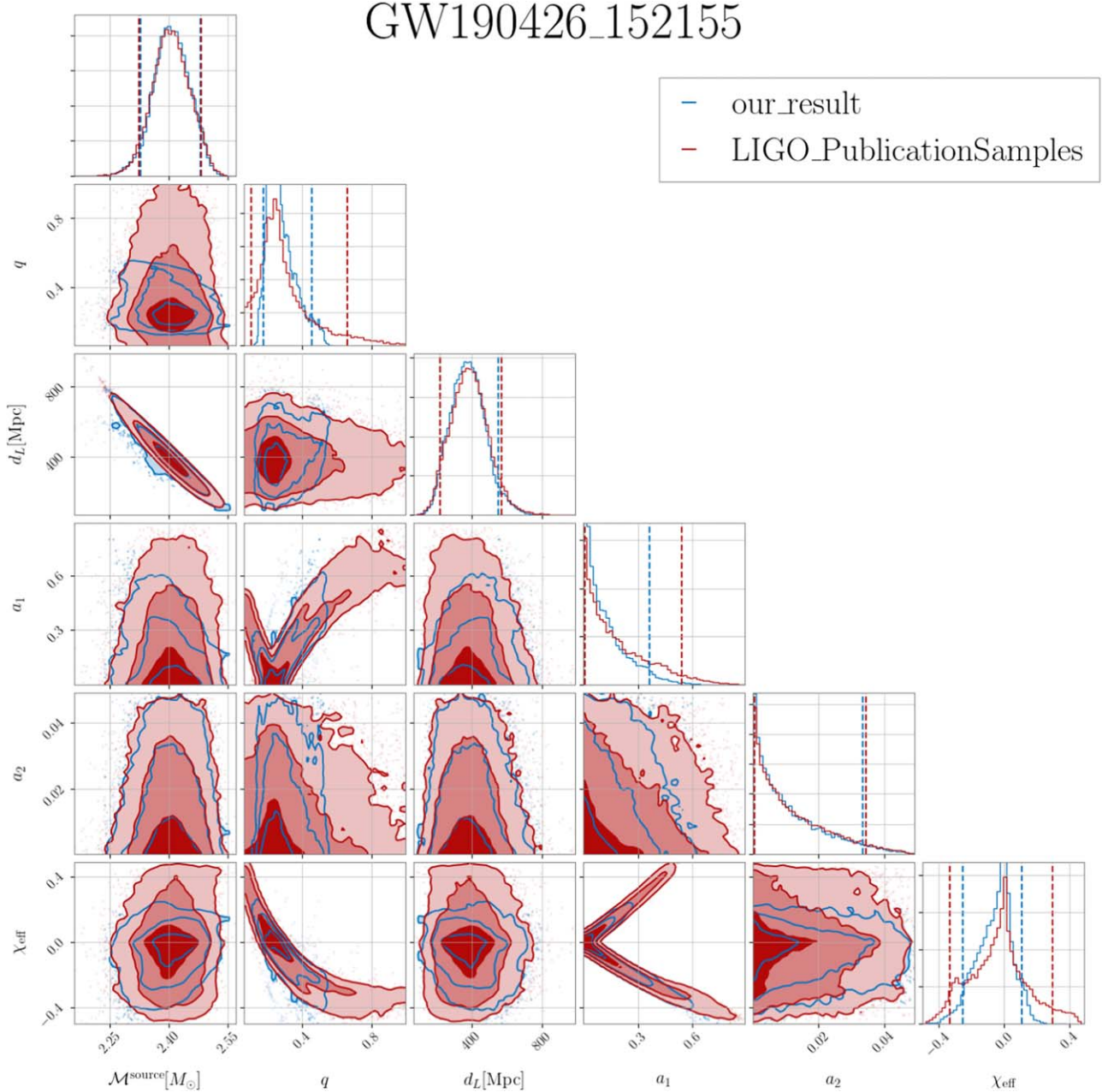


Figure 7. Comparison between our sampling results and posterior samples released by LVC for the event GW190426_152155. The blue regions and lines denote our results and the red for LVC. The dashed vertical lines denote the 5% and 95% quantiles. The labels of parameters follow the conventions in `bilby`. Since we impose a constraint on the prior of the secondary mass, the distribution of mass ratio and effective inspiral spin parameter have slight differences with LVC. Our sampling results are consistent with the results released by LVC within tolerance.

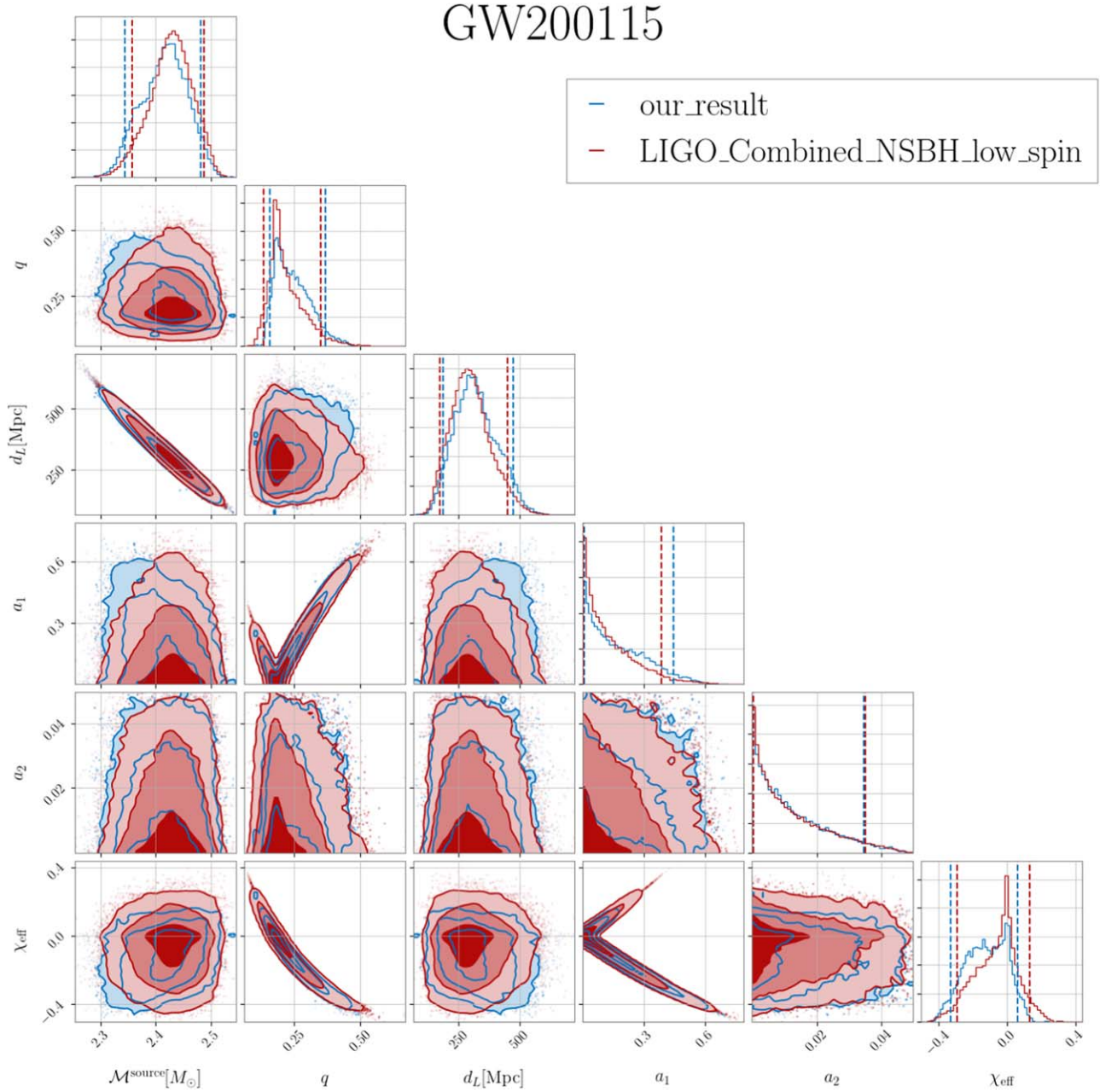


Figure 8. Comparison between our sampling results and posterior samples released by LVC for the event GW200115. Keeping the same with the last figure, the blue and red colors are used to denote our results and LVC's, and dashed vertical lines indicate the intervals of 90% CL. The posteriors of GR parameters of the event GW200115 in our runs are also consistent with the results released by LVC within tolerance.

All results of our parameter estimation can be found on Zenodo.¹⁰ The differences between all our results and LVC's are within tolerance.

Appendix C

Comparison of the Constraints on Dipole Radiation

As discussed in Section 3, it is practically difficult to constrain α_0 or (α_0, β_0) by the events GW170817 and GW190425. However, these two events can be used to constrain the dipole radiation without considering specific model parameters. In order to compare

the results given by LVC, we also perform the tests on φ_{-2} for these two events. We follow the method of model-independent parameterized tests used by LVC (Abbott et al. 2016a, 2019a, 2019b, 2021d), except that we only consider the physical range of $\varphi_{-2} < 0$ which represents the positive outgoing energy flux. Following the works of LVC (Abbott et al. 2019c, 2020a), we use the pre-processed data in which the glitches have been subtracted (Cornish & Littenberg 2015; LIGO Scientific Collaboration & Virgo Collaboration 2017, 2018; Pankow et al. 2018; LIGO Scientific Collaboration & Virgo Collaboration 2019a; Davis et al. 2019; Driggers et al. 2019) and event-specific PSDs encapsulated in LVC posterior sample releases (LIGO Scientific Collaboration & Virgo Collaboration 2019b, 2020b) to perform full Bayesian inference.

⁹ <https://dcc.ligo.org/LIGO-P2100143/public> for GW200115 and <https://dcc.ligo.org/LIGO-P2000223/public> for GW190426_152155.

¹⁰ <https://doi.org/10.5281/zenodo.4817420>

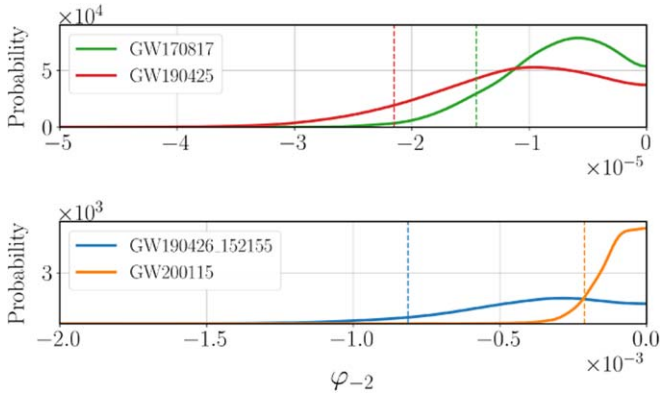


Figure 9. Comparisons between the posterior distributions of φ_{-2} . The dashed vertical lines denote the limits at 90% CL. The limit for GW170817 is about 10^{-5} which is consistent with the result reported by LVC (Abbott et al. 2019b). The limit provided by GW190425 is comparable with GW170817, only have a slight difference within the same order of magnitude. While the limits given by the two NSBH events are much worse than the limits given by the BNS events. The better constraints are because the BNS events have lighter masses which allows more circles of inspiral to be observed in the detectors sensitive band.

The results are shown in Figure 9. The limits at 90% CL are shown by the dashed vertical lines. The limit for GW170817 is about 10^{-5} which is consistent with the result reported by LVC (Abbott et al. 2019b). The limit provided by GW190425 is comparable with GW170817, only have a slight difference within the same order of magnitude. While the limits given by the two NSBH events are much worse than the limits given by the two BNS events. The better constraint is because the BNS events have a lighter mass which allows more circles of inspiral to be observed in the detectors sensitive band. For the same reason, the limit given by GW200115 is slightly better than limit given by GW190426_152155.

Appendix D

Relations between the Scalar Charge and the Mass for Different EoSs

The EoS has to be given in order to solve the TOV equation. Considering the measurements of the millisecond pulsar PSR J0030+0451 and PSR J0740+6620 (Miller et al. 2019, 2021; Riley et al. 2019, 2021) and observational evidence that the maximum mass of neutron stars can exceed $2M_{\odot}$ (Demorest et al. 2010; Antoniadis et al. 2013; Fonseca et al. 2016; Arzoumanian et al. 2018; Cromartie et al. 2019), we select four commonly used

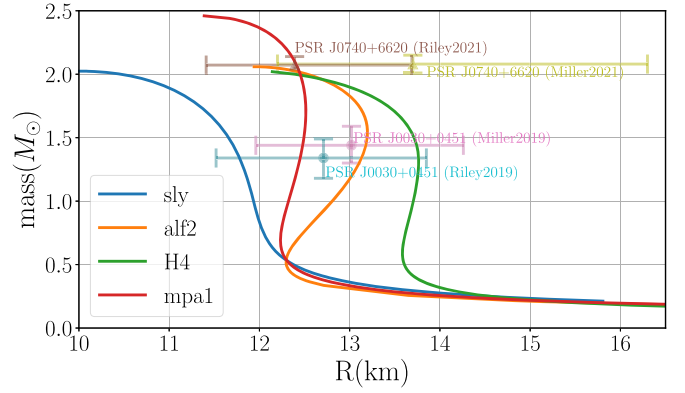


Figure 10. The relation between mass and radius in GR for four EoS used in this work. The four solid lines denote the different EoSs, and the translucent error bars denote the 68% credible regions of mass–radius measurements of the millisecond pulsars PSR J0030+0451 and PSR J0740+6620. The pink one indicates the result reported in Miller et al. (2019) and the cyan is for the result given by Riley et al. (2019). The brown and olive denote the most recent results of PSR J0740+6620 from Riley et al. (2021) and Miller et al. (2021), respectively. We select these four EoSs by considering these observation constraints on mass–radius relation and the observational evidence that the mass of a neutron star can exceed $2M_{\odot}$ (Demorest et al. 2010; Antoniadis et al. 2013; Fonseca et al. 2016; Arzoumanian et al. 2018; Cromartie et al. 2019).

EoSs, *sly*, *alf2*, *H4*, and *mpa1*, in this work. We illustrate the relations between mass and radius in GR of these EoSs¹¹ and the measurements of pulsar mass and radius form two independent groups in Figure 10. The four solid lines represent the EoS used in this work, and the translucent error bars indicate the 68% credible regions of mass–radius measurements.

Using the four EoS, we can solve the TOV equations by the process discussed in Section 2 and extract the scalar charges and mass from the solutions by the Equations (A4). The relations between the mass and the scalar charge are shown in Figure 11 for BD and Figure 12 for DEF. For BD, as can be seen in Figure 11, the influence of using different EoS is slight. The differences of the scalar charge (relative to α_0) are within the order of 0.1. Unlike BD, the curves represent the relation between scalar charge and mass have apparent differences for different EoSs in DEF. The same conclusion also was presented in previous works (Shibata et al. 2014; Shao et al. 2017). For different EoSs, the magnitude of scalar charges amplified by scalarization phenomena is almost same, but the scalarization windows which is the mass range where the nonperturbative strong-field effects can occur are different.

¹¹ Data used to plot are downloaded from http://xtreme.as.arizona.edu/NeutronStars/data/mr_tables.tar.

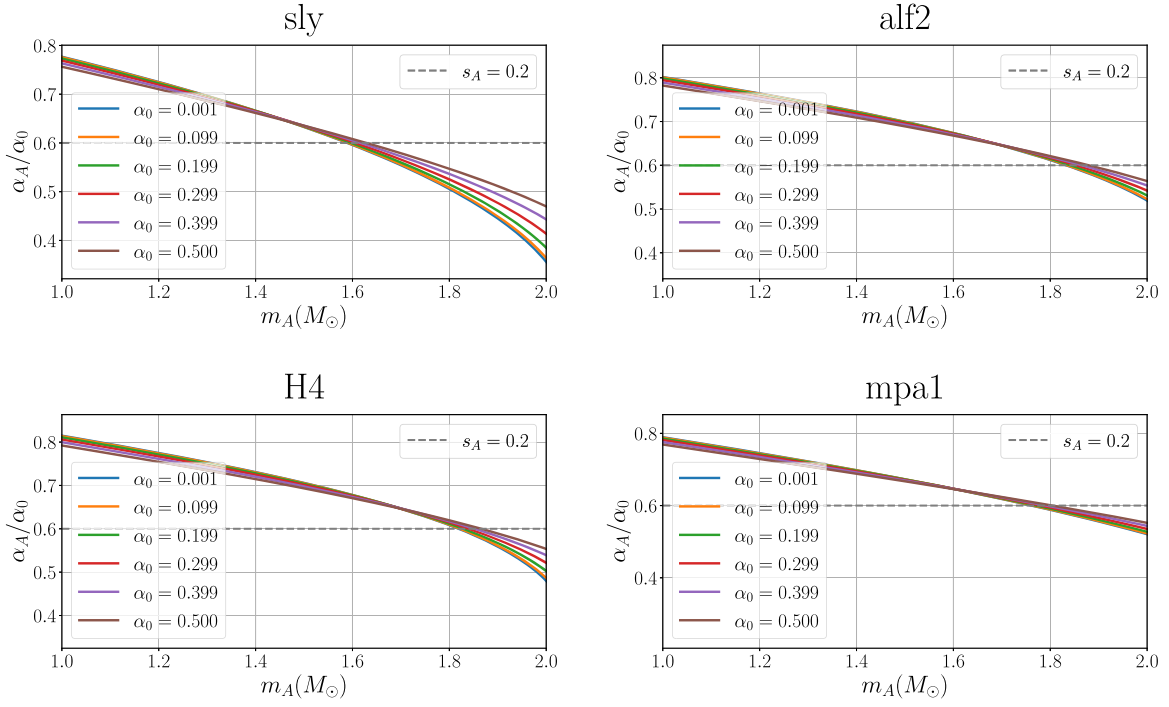


Figure 11. The relation between scalar charge and mass in BD for four different EoSs. Different colors are used to denote different values of α_0 . These results show the relation is similar for different EoSs in BD. The differences of the scalar charge (relative to α_0) are within the order of 0.1.

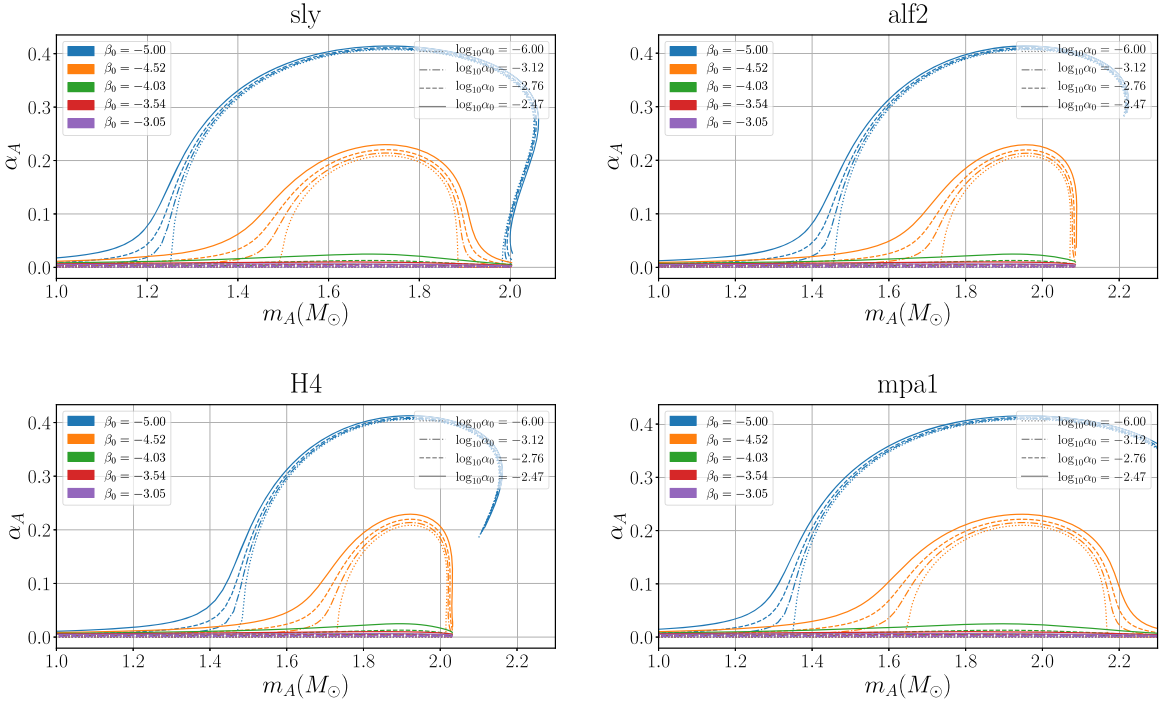


Figure 12. The relation between scalar charge and mass in DEF for four different EoSs. The colors are used to denote the values of β_0 and the line styles are used to denote the values of α_0 . For different EoSs, the magnitude of scalar charges amplified by scalarization phenomena is almost same, but the scalarization windows which is the mass range where the nonperturbative strong-field effects can occur are different.

Appendix E

Unphysical Deviations of Highly Asymmetric Sources

In the main body, we exclude the two events GW190814 and GW200105 due to unphysical deviations. Here, we show the posteriors of the dipole modification parameter for these two events in Figure 13 and present more discussion below.

Subdominant spherical harmonic multipoles will become important when the mass ratio of sources is large. There is strong evidence for the presence of higher modes (HMs) in the analysis of GW190814 (Abbott et al. 2020c). Therefore, following the paper of LVC (Abbott et al. 2021d), we also employ the waveform model incorporating HMs, IMRPhenomPv3HM

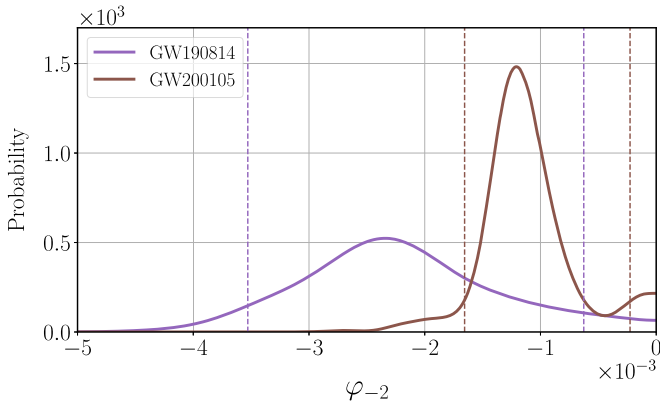


Figure 13. The posteriors of the dipole modification parameter φ_{-2} for the two NSBH events excluded in this work. The dashed vertical lines indicate 5% and 95% percentiles for the two events, respectively. In the posterior of GW190814, the best fit value deviates from the GR value in the order of 10^{-3} , the GR value falls in the tail and is excluded from the 90% confidence interval. The deviation shown here is in agreement with the LVC analysis which can be seen in Figure 19 of Abbott et al. (2021d). The similar deviation is also present in the result of GW200105.

(Khan et al. 2020). The waveform model IMRPhenomPv3HM is based on the model IMRPhenomD (Husa et al. 2016; Khan et al. 2016) which is employed in the main body, but incorporates the processing due to the in-plane spins and HMs (London et al. 2018; Khan et al. 2019, 2020). The same is true for LVC (Abbott et al. 2021d); we only add the dipole modification on the dominant mode. The non-GR deformation on HMs is obtained by rescaling the modification in the dominant mode according to the method presented in London et al. (2018). There are no new coefficients introduced. It is worth noting that this method of implementation can possibly be one of the reasons that causes unphysical deviation.

Using this waveform model, we perform the same Bayesian inference discussed in the main body on the two events GW190814 and GW200105 to constrain the dipole modification parameter φ_{-2} . The results are shown in Figure 13. The dashed vertical lines indicate 5% and 95% percentiles for the two events respectively. It can be seen that the GR value falls in the tails of the posteriors and is excluded from the intervals of 90% CL for these two events. The best fit value of GW190814 deviates from the GR value in the order of 10^{-3} . While the deviation of GW200105 is slightly reduced. The result of GW190814 presented here is consistent with the result of LVC (as can be seen in Figure 19 of Abbott et al. 2021d). Similar deviations are also reported in Perkins et al. (2021). These results are believed to be not the real deviations from GR. The possible reasons for these deviations might be the systematic errors of the waveform templates or the parameterization method of non-GR modification (which might be inappropriate when HMs are present as pointed out above), and covariances between model parameters (Abbott et al. 2021d; Perkins et al. 2021).

We also find the deviations are somehow related to the mass parameters of sources. Referring to Figure 4 in Abbott et al. (2021a), we also illustrate the component masses of all four possible NSBH events so far in Figure 14. The 90% CL regions of the joint posterior distribution for component masses are enclosed by the solid curves, and the shading denotes the posterior probability. The dashed gray lines indicate the constant mass ratio.

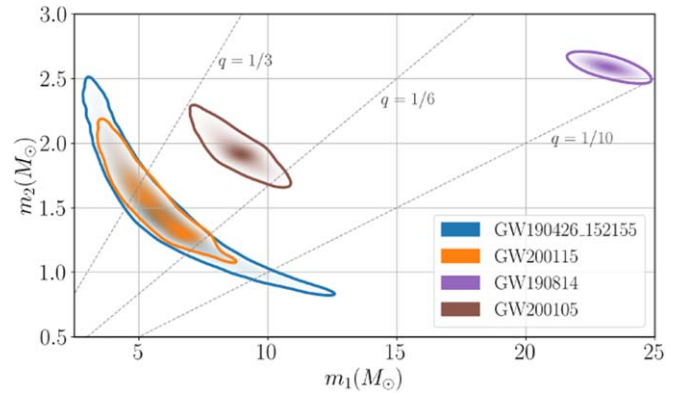


Figure 14. The component masses of all four possible NSBH events so far. Following Figure 4 in Abbott et al. (2021a), we illustrate the component masses of the four possible NSBH events for convenience of reference. The 90% CL regions of the joint posterior distribution for component masses are enclosed by the solid curves, and the shading denotes the posterior probability. The dashed gray lines indicate the constant mass ratio. As can be seen in the figure, the GW190814 is the most asymmetric source. Deviations that might be caused by systematic errors of waveform templates, the parameterization method of non-GR modification, or covariances between model parameters, are present in the posteriors of the dipole modification parameter as shown in Figure 19 in Abbott et al. (2021d), Figure 10 in Perkins et al. (2021), and Figure 13 in this paper. Similar deviations are also seen in the case of GW200105, but absent in GW190426_152155 and GW200115, which is probably due to the more dispersed posteriors and greater probability on the lower mass ratio.

The posterior distributions of GW190426_152155 and GW200115 are almost overlapping. The posteriors of these two events are the most dispersed and have more in the lower mass ratio. Meanwhile, the deviations on φ_{-2} are absent for these two events. The events GW190814 and GW200105 have higher mass ratio, and the magnitude of deviation from GR value is consistent with their mass ratio as can be observed by combining Figures 14 and 13.

For sources with large mass ratios, the HMs become more important which may complicate the analysis. As discussed above, the non-GR modifications on HMs are propagated from the rescaling of the modifications on the dominant mode according to the rules presented in London et al. (2018). The rescaling rules are verified for the GR part by numerical relativity but are doubtful for the non-GR part. The method of performing the parameterized tests may be inapplicable when HMs are present. The parameterized tests have not been systematically studied in the parameter space of highly asymmetric sources. More thorough studies are needed to explain these deviations. In this work, we simply exclude the two events GW190814 and GW200105.

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