

Direct Multimodel Dark Matter Search with Gravitational-Wave Interferometers Using Data from the First Part of the Fourth LIGO-Virgo-KAGRA Observing Run

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We present a search of data from the first part of the fourth observing run of LIGO-Virgo-KAGRA for three kinds of dark matter—dilaton (spin-0), dark photons (spin-1), and tensor bosons (spin-2)—using three independent methods, each extended to search for all three candidates. Each dark matter candidate could interact with different standard-model particles in the instruments, causing unique differential strains on the interferometers. While we do not find any evidence for a signal, we place the most stringent upper limits to date on the couplings of each dark matter model to the interferometers. For scalars with masses between $[4 \times 10^{-14}, 1.5 \times 10^{-13}]$ eV that couple to photons or electrons, our constraints on the scalar-to-electron or photon coupling improve upon those from the third observing run by 1 order of magnitude, with the tightest limit of $\sim 10^{-20} \text{ GeV}^{-1}$ at a mass of $\sim 2 \times 10^{-13}$ eV. For vectors with masses between $[7 \times 10^{-13}, 8.47 \times 10^{-12}]$ eV that couple to baryons, our constraints on the dark-photon or baryon coupling supersede those from MICROSCOPE and Eöt-Wash by 1 to 2 orders of magnitude, reaching a minimum of $\sim 5 \times 10^{-24}$ at a mass of $\sim 10^{-12}$ eV. For tensors with masses of $[4 \times 10^{-14}, 8.47 \times 10^{-12}]$ eV (the full mass range analyzed) that couple via a Yukawa interaction, our constraints, which are the first obtained from gravitational-wave interferometers, surpass those from fifth-force experiments by 4 to 5 orders of magnitude, achieving a limit as low as $\sim 8 \times 10^{-9}$ at $\sim 2 \times 10^{-13}$ eV. Our results show that each method has substantially different sensitivity to each dark matter candidate, highlighting the importance of a multimethod approach to search for new physics. Moreover, they demonstrate that gravitational-wave interferometers can simultaneously probe multiple ultralight dark matter models and place unprecedented constraints on tensor fields arising in a broad range of modified-gravity theories.

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Introduction—The existence of dark matter (DM) is well established, yet its fundamental nature remains one of the most pressing mysteries in modern physics [1]. Across a vast range of possible masses and interaction strengths, DM could take many forms—from particles with masses of $\mathcal{O}(10^{-22})$ eV, whose wavelike behavior spans galactic scales [2], to primordial black holes with masses comparable to stars [3–5]. Among these possibilities, ultralight DM with masses $m_{\text{DM}} \ll 1$ eV presents a particularly compelling scenario. Not only do high-energy theories naturally predict such particles, but their extraordinarily low mass also implies a quantum-mechanical coherence over macroscopic distances, effectively turning them into oscillating background fields. If these fields interact weakly with standard-model particles, they could produce

distinctive, nearly monochromatic signals detectable by precision experiments [6,7].

Gravitational-wave (GW) detectors such as LIGO, Virgo, and KAGRA [8–10] have already reshaped astrophysics by measuring the signals from merging black holes and neutron stars [11–13]. But their extraordinary sensitivity also makes them powerful tools to probe fundamental physics beyond GWs [14–17]. The same precision that captures the faint signals from GWs could reveal subtle interactions between DM and ordinary matter, and probe small modifications to gravity, which places these detectors at the forefront of searches for new physics. Along these lines, stringent constraints have been set on dilatons and dark photons from searches on past GW data [18–23].

The ultralight DM candidates considered here occupy a portion of the parameter space that is largely inaccessible to previous experiments. Indirect-DM searches via gamma-ray emission, e.g., VERITAS or Fermi, target weakly interacting massive particles in the GeV–TeV mass range [24], while experiments such as NuSTAR constrain sterile-neutrino DM at masses of $\mathcal{O}(\text{keV})$ [25]. Additionally, ADMX [26] and ALPS [27] have searched for axion ultralight DM with masses of $\mathcal{O}(\mu\text{eV})$ and $\mathcal{O}(\text{meV})$,

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respectively, which also lie outside the mass range that we study.

Likewise, existing GW searches with LIGO, Virgo, and KAGRA are not sensitive to the ultralight DM models considered here, despite targeting signals with similar morphologies. These searches are designed to find Doppler-shifted, spinning-down signals from rapidly rotating neutron stars or boson clouds around spinning black holes [28], which do not match the signatures of ultralight DM signals. Additionally, narrow-band stochastic GW background searches use a frequency-dependent overlap reduction function (ORF) that differs from that of ultralight DM interactions [29,30] [31], and much shorter analysis coherence times that are not optimized for ultralight DM. Moreover, these searches only consider real, positive signal-to-noise ratios (SNRs) as being astrophysical, whereas both dark photons and tensor DM would generate negative SNRs because of their ORFs [32], causing standard analyses to miss them [29,30].

In contrast to previous works, we explore here how GW detectors can be used to search for ultralight DM in various forms simultaneously: scalar fields (such as dilatons, which may induce oscillations in fundamental constants) [7,33–35], vector fields (like dark photons, exerting weak oscillating forces on matter) [6], and massive tensor fields (closely resembling GWs but with a nonzero mass) [36–39], which appear in a broad range of modified gravity theories. Each of these candidates leaves a unique imprint on GW detectors, offering a new way to constrain—or perhaps discover—some of the most elusive DM candidates predicted by theory.

Dark matter interaction models—Cold, ultralight DM consists of an enormous number of extremely light particles, which motivates its treatment as a classical oscillating field. The frequency of this field is fixed by the ultralight DM mass regardless of the type of DM considered: $f_0 = m_{\text{DM}} c^2 / h$; however, there is a small, stochastic frequency change caused by the Maxwell-Boltzmann velocity distribution of the DM particles: $\Delta f / f \sim v_0^2 / c^2 \sim 10^{-6}$, where v_0 is the virial velocity of DM around the center of the Galaxy. This frequency spread cannot be captured by standard deterministic methods that look for quasimonochromatic signals without incurring a sensitivity loss. Moreover, the signal has a finite coherence time that scales inversely with its mass, of $\mathcal{O}(10^4)$ s at $m_{\text{DM}} \sim 10^{-12}$ eV [17], in which the signal can be treated as purely monochromatic, but spreads over Δf if observing for longer than the coherence time. Likewise, the signal has a finite coherence length, which greatly exceeds the separation of Earth-based interferometers, implying correlated signals.

Regardless of the ultralight DM model, the signal appears in the detector as a persistent, quasimonochromatic signal with a finite coherence time, and a narrow-band signal with stochastic frequency fluctuations over longer timescales. Each model described below for spin-0, 1, and 2

differs in the mechanism by which DM couples to the interferometers and in the induced strain amplitude, not in the phase evolution. While models for higher-spin DM exist [40–42], the ways in which they interact with GW interferometers are model-dependent. After specifying the couplings, we can generalize the analyses presented here to constrain them as well.

Scalar, dilaton dark matter: Models of ultralight DM predict that it could have originated through a vacuum misalignment mechanism in the early Universe, and would manifest today as a coherently oscillating scalar field [34]. The coupling of such DM would effectively induce an oscillation of the values of fundamental constants, namely the electron rest mass and the fine structure constant, thus leading to the expansion and contraction of solids. When a field gradient exists, objects can also be accelerated. Many ideas have been proposed to detect such effects [43–46], with the current strongest direct constraints coming from experiments looking for the oscillation of fundamental constants using atomic clocks or spectroscopy [47–58]. More recently, several experimental mechanisms, in particular the way in which dilatons could couple to the beam splitter, were discovered to detect those same oscillations using GW interferometers [7,33–35,59,60].

Considering a simple linear expansion to the standard-model Lagrangian $\mathcal{L}$, we write the scalar field ϕ as [20]

$$\mathcal{L} \supset \frac{\phi}{\Lambda_\gamma} \frac{F_{\mu\nu} F^{\mu\nu}}{4} - \frac{\phi}{\Lambda_e} m_e \bar{\psi}_e \psi_e, \quad (1)$$

where $F_{\mu\nu}$ is the electromagnetic field tensor, ψ_e and $\bar{\psi}_e$ are the standard-model electron field and its Dirac conjugate, respectively, m_e is the electron rest mass, and Λ_γ and Λ_e denote the scalar DM coupling constants to the photon and electron, respectively. When Λ_e^{-1} is nonzero, the DM field induces changes in m_e and the fine structure constant through $F_{\mu\nu}$. In turn, this modulates the sizes and refractive indices of solids.

In LIGO, by considering only oscillations in the size of the beam splitter and mirrors, we can relate the measured strain $h(f_0)$ to Λ_i as [23]

$$\begin{aligned} \left(\frac{1}{\Lambda_\gamma} + \frac{1}{\Lambda_e} \right) &\approx \frac{m_{\text{DM}} c}{\hbar \sqrt{2 \rho_{\text{DM}}}} h(f_0) A_{\text{cal}}(f_0) \\ &\approx 10^{-18} \text{ GeV}^{-1} \left(\frac{f_0}{100 \text{ Hz}} \right) \left(\frac{h_0(f_0) A_{\text{cal}}(f_0)}{10^{-17} \text{ s}^{-1}} \right), \end{aligned} \quad (2)$$

where $\rho_{\text{DM}} = 0.4 \text{ GeV cm}^{-3}$ is the local DM energy density from the standard galactic DM halo model [61]. A_{cal} encodes the interferometer response to scalar DM [23], and is obtained through optical simulations that account for finite light travel time effects [60] (see Supplemental Material [62]). Note here we make the minimal assumption about the

scalar potential by keeping only the quadratic mass term. This type of potential can naturally arise when the presence of an underlying approximate symmetry leads to strongly suppressed self-interactions, such as the approximate shift symmetry in relaxion models [63]. In such models, the quartic and higher-order terms are small and do not induce a significant change in the scalar mass near the Earth due to the scalar field's linear couplings to matter [64–68].

Vector dark-photon dark matter: Spin-1 dark photons [69] could completely explain the relic abundance of DM [71]. These particles could arise from the misalignment mechanism [72–74], parametric resonance or the tachyonic instability of a scalar field [75–78], or from cosmic string network decays [79]. The observable effect would result from a coupling of dark photons to standard-model particles—either baryons or baryon minus leptons. In particular, this interaction would cause “dark” forces on the mirrors that would oscillate at a frequency fixed by dark photon [6,80].

The Lagrangian $\mathcal{L}$ that characterizes the dark-photon coupling to a number current density J^μ of baryons or baryons minus leptons is

$$\mathcal{L} = -\frac{1}{4\mu_0} F'^{\mu\nu} F_{\mu\nu} + \frac{1}{2\mu_0} \left(\frac{m_{\text{DM}C}}{\hbar}\right)^2 A^\mu A_\mu - e e J^\mu A_\mu, \quad (3)$$

where $F_{\mu\nu}$ is the dark-photon electromagnetic field tensor, μ_0 is the magnetic permeability in vacuum, e is the electron charge, A_μ is the dark four-vector potential, and ϵ is the strength of the particle or dark-photon coupling normalized by the electromagnetic coupling constant.

The strain on the interferometers caused by a dark photon DM signal is [60]

$$\begin{aligned} \sqrt{\langle h_D^2 \rangle} &= \frac{1}{3\sqrt{2}\pi} \frac{Q}{M} \frac{\hbar e}{c^4 \sqrt{\epsilon_0}} \sqrt{2\rho_{\text{DM}} v_0} \frac{\epsilon}{f_0}, \\ &\simeq 6.56 \times 10^{-26} \left(\frac{\epsilon}{10^{-22}}\right) \left(\frac{100 \text{ Hz}}{f_0}\right), \end{aligned} \quad (4)$$

where ϵ_0 is the permittivity of free space, and Q/M is the dark charge-to-mass ratio of the LIGO mirrors under the new $U(1)_B$ current. We take $Q/M = m_p^{-1}$ for $U(1)_B$ because the mass of the mirror is dominated by protons and neutrons [81], where m_p is the proton rest mass [82].

A second effective arm-length modulation results because light takes a finite amount of time to travel between the input and end mirrors in the interferometer, during which the mirrors have moved in response to DM [60],

$$\begin{aligned} \sqrt{\langle h_C^2 \rangle} &= \frac{\sqrt{3}}{2} \sqrt{\langle h_D^2 \rangle} \frac{2\pi f_0 L}{v_0}, \\ &\simeq 6.58 \times 10^{-25} \left(\frac{\epsilon}{10^{-22}}\right), \end{aligned} \quad (5)$$

where $L = 4 \text{ km}$ is the arm length of the LIGO interferometers. The total strain from both physical effects always adds as $\langle h_{\text{total}}^2 \rangle = \langle h_D^2 \rangle + \langle h_C^2 \rangle$. The averages are taken over polarization direction and detector geometry.

Tensor dark matter: A novel spin-2 graviton arising in ghost-free massive gravity theories can simultaneously modify gravity and constitute DM [83–86]. Scenarios in which massive gravitons comprise all of DM have been widely explored [32,36,37,42,84,87–90], along with the associated phenomenology and detection methodologies in ground-based or space-based laser interferometers [32,38,39,91], atom interferometers [92], and pulsar timing arrays [93–97], and through black-hole superradiant instabilities [98]. In this Letter, we consider spin-2 tensor DM arising from the Hassan-Rosen bigravity [85], consisting of one massless and one massive graviton, where the massive eigenstate with 5 degrees of freedom could explain DM [37].

The Fierz-Pauli Lagrangian density [95] describes this tensor field $\chi_{\mu\nu}$,

$$\begin{aligned} \mathcal{L} &= -\frac{1}{4} \chi^{\mu\nu} \mathcal{E}_{\mu\nu,\alpha\beta} \chi^{\alpha\beta} - \frac{1}{8} \left(\frac{m_{\text{DM}C}}{\hbar}\right)^2 (\chi_{\mu\nu} \chi^{\mu\nu} - \chi^2) \\ &\quad - \frac{\alpha}{2cm_{\text{Pl}}} \sqrt{\frac{\hbar}{c}} \chi_{\mu\nu} T^{\mu\nu}, \end{aligned} \quad (6)$$

where m_{Pl} is the reduced Planck mass, $\chi \equiv \chi^\mu_\mu$, $\mathcal{E}_{\mu\nu,\alpha\beta}$ is a second-order derivative operator defined in Eq. (2.13) of [36], and α is the Yukawa coupling to the energy-momentum tensor $T_{\mu\nu}$. This kind of coupling is generic to ultralight tensor DM models [36].

The strain on interferometers arises analogously to that from GWs, a stretching of space-time in the presence of the field. Specifically, ignoring the finite-light-travel-time effect, the differential strain after averaging over five polarization components can be expressed as [36,39,99]

$$\begin{aligned} h(f_0) &= \frac{2\alpha}{2\pi f_0 m_{\text{Pl}}} \sqrt{\frac{\rho_{\text{DM}} \hbar c}{5}} \\ &\simeq 7.78 \times 10^{-26} \left(\frac{\alpha}{5 \times 10^{-8}}\right) \left(\frac{100 \text{ Hz}}{f_0}\right). \end{aligned} \quad (7)$$

Data—We analyze data from the first part of the fourth observing run (O4a) of the LIGO Livingston (L1) and Hanford (H1) detectors. The sensitivity of LIGO in O4a improved significantly across all frequencies, particularly above 400 Hz [100–102]. The strain data were collected between May 24, 2023, at 15:00 UTC, and January 16, 2024, at 16:00 UTC, with L1 and H1 operating 69.0% and 67.5% of the time, respectively [103]. Virgo did not take data during O4a, and KAGRA observed for one month. The data are calibrated by directing a laser onto the test masses and comparing the measured response to the expected one [104,105]. Additionally, the data have been cleaned to remove “glitches,” short-duration disturbances

[106], and analyzed only if they are in “science mode” [107]. At worst, amplitude and phase uncertainties at 1σ are 10% and 10° , respectively, in the frequency range analyzed by all methods: [10, 2000] Hz.

In this analysis, we use three independent pipelines, described in “Search methods,” to search for direct DM interactions with GW interferometers. Each one employs data in different formats. The first, called the “BSD-excess-power” method, uses band sampled data (BSD) structures [108] as inputs to this search, which are derived from short fast Fourier transform databases [109]. BSDs consist of complex-valued time-series representations of the strain data in 10-Hz bands, generated monthly from the GDS-CALIB_STRAIN_CLEAN_AR channel. The other two methods, “cross-correlation” and logarithmic power spectral density (LPSD), use *gated* data [110] from the channel GDS-CALIB_STRAIN_CLEAN_GATED_G02, in which spurious glitches have been removed [111,112]. Cross-correlation employs short Fourier transforms (SFTs) of length $T_{\text{SFT}} = 1800$, totalling 5696 coincident H1–L1 SFTs (2848 h). In contrast, LPSD requires continuous segments of at least 10^5 s, resulting in 30 coincident segments (1157 h).

Search methods—We present three independent methods that searched for ultralight DM coupling to the interferometers. All methods are sensitive to each kind of DM presented in the introduction because each only looks for a monochromatic signal, which is a shared feature of ultralight DM signals coupling to GW interferometers. The first, cross-correlation, relies more heavily on the detector response to each kind of DM and is therefore inherently model-dependent. In contrast, BSD excess power and LPSD are generically sensitive to narrow-band ultralight DM, allowing null results to be more easily reinterpreted as constraints on different DM couplings. For cross-correlation, extending the search to different kinds of DM required model-specific adaptations to the search, including the derivation of new ORFs. In contrast, for BSD excess power and LPSD, the search strategies are unchanged relative to previous works. For all three methods, the resulting strain upper limits are translated into constraints on all three DM couplings for the first time.

Cross-correlation: Cross-correlation is a natural technique to search for ultralight DM because the signal is expected to correlated among multiple detectors [6], and was used in O1 and O3 DM searches [18,19]. By cross-correlating the data in each frequency bin j , we can calculate the SNR as

$$\text{SNR}_j = \frac{S_j}{\sigma_j}. \quad (8)$$

S_j is the signal strength constructed by cross-correlating fast Fourier transform (FFT) datasets $z_{1,ij}$ and $z_{2,ij}$ of detectors 1 and 2 for segment i , and then averaging over the segments,

$$S_j = \frac{1}{N_{\text{SFT}}} \sum_{i=1}^{N_{\text{SFT}}} \frac{z_{1,ij} z_{2,ij}^*}{P_{1,ij} P_{2,ij}}, \quad (9)$$

where N_{SFT} is the number of SFTs used, $P_{1(2),ij}$ is the noise power estimated from the neighboring 50 frequency bins [113] using a running median, and is related to the one-sided power spectral density (PSD) by $\text{PSD}_{1(2),ij} = 2P_{1(2),ij}/T_{\text{SFT}}$. When a signal is not present, the expectation value of S_j is zero, and the variance for its real and imaginary parts is

$$\sigma_j^2 = \frac{1}{N_{\text{SFT}}} \left\langle \frac{1}{2P_{1,ij} P_{2,ij}} \right\rangle_{N_{\text{SFT}}}. \quad (10)$$

For scalar DM, the SNR is positive, whereas for dark photons and tensor bosons, it is negative because the corresponding ORF is ~ -0.9 over the analyzed frequency range. These ORFs are derived in Supplemental Material [62].

BSD-excess-power method: As done in O3 [19], in each 1-Hz band, BSD excess power changes the fast Fourier transform coherence time T_{FFT} to match the DM signal coherence time T_{coh} , ensuring that all power is confined to one frequency bin for the full observing run for each DM mass [80]. We then take Fourier transforms of the strain data using $T_{\text{FFT}} \in [10^{2.75}, 10^{4.5}]$ s at high and low frequencies, respectively [19,80], to make time-frequency “peak maps” [109,114]. Peak maps are collections of ones that contain frequency bins whose powers exceed a threshold in the equalized spectrum, i.e., the ratio of the signal power normalized by the noise power, estimated with a running median [80], and are also local maxima. We then count the number of peaks at each frequency to find bins with large numbers of peaks. By counting the peaks, and not the power in each frequency bin, we are more robust against non-Gaussian noise artifacts that could contaminate particular frequency bins and bias our detection statistic, the critical ratio CR,

$$\text{CR} = \frac{y - \mu}{\sigma}. \quad (11)$$

Here, y is the number of peaks at a particular frequency, and μ and σ are the mean and standard deviations of the number counts across all frequencies in the band, respectively. In Gaussian noise, the CR is normally distributed with zero mean and unit variance [114], and a signal would appear in the positive tail of this distribution.

We select enough candidates [115] in each 1-Hz band such that one coincident candidate between two detectors would occur in Gaussian noise. This ensures that we select uniformly over the frequency range.

LPSD: LPSD exactly matches the integration time in Fourier transforms to the DM coherence time, similarly to what BSD excess power does, but in every logarithmically spaced frequency bin, not every hertz. This allows it to

achieve optimal sensitivity at each DM mass, regardless of the type of DM. However, the logarithmic frequency spacing poses a challenge because FFT-based algorithms are no longer applicable. With a computational complexity scaling as $\mathcal{O}(N^2)$, where N is the number of data points, this leads to intractable costs. However, we developed a method that leverages symmetries between the time and frequency domains [116]. This allows us to perform the full logarithmic calculation with a single FFT followed by a heavily zero-suppressed transformation without the need for expensive precomputations.

We then follow the approach in [23,116], where a background model is constructed by fitting cubic splines in log-log space to the PSD as a function of frequency. Because the PSDs vary by several orders of magnitude, we fit over intervals of 10^4 frequency bins. Within these intervals, potential DM signals are identified as positive deviations from the fits by using a profile likelihood test [see Eq. (15) in [116]], assuming, as validated empirically, that the residuals follow a skew-normal distribution.

Results—

Cross-correlation: We set thresholds for detection of $\text{Re}(\text{SNR}) < -5.8$ for dark photons and tensor bosons, and $\text{Re}(\text{SNR}) > 5.8$ for scalar bosons that correspond to $\sim 1\%$ false-alarm probability, but find no candidates. We also checked subthreshold outliers with SNRs in the range [5.0, 5.8] for either the real or imaginary part of the SNR. Of ten such outliers, nine were due to loud instrumental artifacts, and the remaining one was consistent with Gaussian noise.

In the absence of a detection, and after excluding frequency bins lying within ~ 0.056 Hz of known spectral artifacts, the measured SNR is converted into upper limits on each of the aforementioned DM couplings using the Feldman-Cousins formalism [117]. This conversion incorporates both an efficiency factor obtained from simulations to correct for the power loss due to the frequency binning choice and random frequency-bin boundary, and the appropriate ORF. These limits are set bin by bin and thus exhibit larger fluctuations than limits calculated in 1-Hz bands by the other methods.

BSD-excess-power method: Our search returned 6847 coincident candidates present in the same frequency bin in both L1 and H1 and that did not overlap with known noise disturbances [106]. However, after requiring that the average CR be greater than 5 (corresponding to 5σ significance in Gaussian noise), only ten candidates remained. To determine whether these candidates resulted from DM, we performed three follow-up procedures: (1) we varied the observation time T_{obs} at the chosen T_{FFT} to see if there was a steady increase in the $\text{CR} \propto T_{\text{obs}}^{1/2}$, (2) we increased $T_{\text{FFT}} > T_{\text{coh}}$ to see if there was a *decrease* in the CR [20,80], and (3) we correlated the spectra from L1 and H1 with an expected DM waveform using the Wiener filter [118] and computed the residual between them [119]. No candidate survived all three follow-ups.

In the absence of candidates, we set Feldman-Cousins [117] upper limits on the minimum detectable strain amplitude at 95% confidence $h_{0,\text{min}}^{95\%}$ induced by a generic DM signal using Eq. (75) of [17], which depends on the detector power spectral density, the maximum of the coincident candidates' CRs, T_{FFT} , and T_{obs} . Up to an $\mathcal{O}(1)$ factor that depends on the DM polarization and the geometry of LIGO, these limits can then be mapped to *any* kind of DM interaction that would induce a differential strain on the detector.

LPSD: We obtain candidates by performing a profile likelihood ratio test in each bin, as outlined in [116], and require a 5σ significance corrected for the “look-elsewhere effect” in rejecting the no-DM hypothesis. After clustering neighboring bins, requiring significance and compatibility in both detectors, we find only one candidate that was found to originate from a nonoptimal setting in the background fit.

The resulting 95% confidence-level upper limits on strain amplitude can then be derived. The use of the gating procedure led to a factor of ~ 3 improvement in the LPSD upper limits with respect to those in O3 [23]. Injection tests were performed that confirmed the upper limits at the expected level, but found a systematic 13% overestimate of the limits. This is small compared to the 30% statistical uncertainty.

Upper limits: In Fig. 1, we show upper limits for each of the three DM models from each pipeline, as argued in [17,120]. The semianalytic way that these limits are computed is detailed in [17], and holds even though the data are not Gaussian, as justified by injection studies [18,23,80]. For all three DM models, we see significant improvements in the constraints on the coupling constants with respect to existing ones. For dilatons, cross-correlation is consistent with or outperforms BSD excess power and LPSD because all three methods are equally sensitive to the finite light-travel time effect. For dark photons, cross-correlation obtains less stringent constraints than those from BSD excess power and LPSD because it is much less sensitive to the finite-light travel time effect [39,60]. For tensor bosons, the upper limits obtained by our pipelines are consistent and significantly outperform fifth-force experiments by several orders of magnitude.

We note that the upper limits resulting from tests of the equivalence principle are stronger than those obtained from our search for scalar DM. In this case, our sensitivity is driven by size modulations of the beam splitter. In laser interferometers, light travels back and forth $\mathcal{O}(10^2)$ times in the Fabry-Perot cavities per encounter with the beam splitter, which significantly reduces our sensitivity. Because of the difference in test-mass thicknesses [23], and the better low-frequency sensitivity and longer observing time, we obtain a substantial improvement below 200 Hz compared to the GEO600 [20] search, which did not suffer from Fabry-Perot suppression.

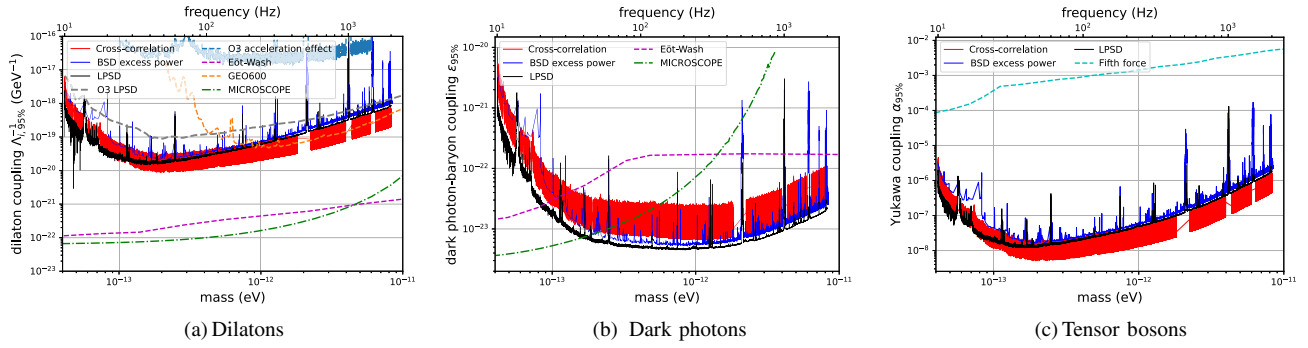


FIG. 1. Upper limits on the coupling strengths of ultralight scalar, vector, and tensor bosons to standard-model particles in LIGO. Solid lines denote the constraints from this work (“cross-correlation,” “BSD excess power,” and “LPSD”). The dashed and dotted-dashed lines indicate constraints from other experiments (“Eöt-Wash” and “MICROSCOPE”) and from previous searches on GW data (“GEO600,” “O3 LPSD,” [We note that the “O3 LPSD curve” is multiplied by a factor of 2.226 compared to the results in [23] after we found a conversion error.] and “O3 acceleration effect”). (a) Our constraints apply to the coupling of dilatons to both photons and to electrons. They improve upon those from previous searches with GW interferometers [20,22,23] and beat those from O3 by one order of magnitude. Constraints from Eöt-Wash [121,122] and MICROSCOPE [123,124] correspond to the coupling of the dilaton to the electron. (b) The constraints on vector dark photon DM improve upon existing limits by a couple orders of magnitudes in the dark-photon and baryon coupling constant. To obtain limits on dark photon and baryon minus lepton coupling, $U(1)_{B-L}$, these limits should be multiplied by approximately two, but depends on the particular mirror composition. (c) For the first time, we derive constraints on tensor DM using GW interferometers, which are up to 5 orders of magnitude superior to those obtained indirectly from fifth-force experiments [125,126]. We release the data to make these figures as well [127].

By contrast, the sensitivity to dark photons and tensor bosons depends primarily on the interferometer arm length, just as for GWs: longer arms imply that smaller displacements of induced DM signals can be measured. In these cases, the differential strain induced by DM coupling to the test masses, or spacetime distortions from tensor fields, is amplified by the long arms. In particular, for tensor bosons, a torsion balance responds as a single bulk object because the coherence length of the DM field is much larger than the size of the experiment [76], leading to an extremely small strain. This is analogous to resonant-bar detectors, which are far less sensitive than interferometers [128]. Thus, fifth-force experiments are stronger probes of scalars than of vectors and tensors, while GW interferometers are more powerful probes of vectors and tensors than scalars.

Conclusions—We have shown that GW interferometers can probe multiple kinds of DM, including tensor bosons arising from a broad range of modified theories of gravity, using three independent search techniques. In particular, we have demonstrated that each of our methods, previously tailored to search for only one particular DM model, can be extended to probe a wide range of DM interactions, enabling a systematic comparison of their sensitivities to scalar, vector, and tensor signals. Cross-correlation is tied explicitly to the signal model through the ORFs and detector response, derived in Supplemental Material [62], and can be used to distinguish among ultralight DM models [39,119]. In contrast, BSD excess power and LPSD are sensitive to generic monochromatic narrow-band signals, independent of the underlying ultralight DM model. This allows null results to be reinterpreted as

constraints on the three DM models considered here, which is done for the first time. Taken together, these three methods combine model-dependent and model-agnostic approaches to search for ultralight DM.

Moreover, we found that each method exhibits substantially different sensitivities to each kind of DM, highlighting the importance of developing multiple independent search pipelines to probe new physics with GW detectors. Specifically, for dilatons and dark photons, detector improvements and increasingly sensitive methods allow us to probe progressively weaker couplings. For tensor bosons, GW interferometers now surpass fifth-force constraints by up to 5 orders of magnitude, establishing them as competitive probes of spin-2 DM and modified gravity theories.

Our chosen tensor-field model serves as a benchmark for a broad range of modified gravity theories, paving the way for our results to be reinterpreted to constrain other tensor fields. We have thus demonstrated that LIGO, Virgo, and KAGRA can act not only as direct DM detectors, but also as direct probes of modified theories of gravity in previously unexplored ways.

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Data availability—The data that support the findings of this article are openly available [127].

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