To compute the integrals in Eq. (26) we adopt the numerical method described in Appendix B. As observed in [50], the integration over the strain amplitude is performed up to h^* to exclude the individually resolvable powerful and rare bursts. The maximum strain amplitude h^* is determined by solving the equation

$$\int_{h^*}^{+\infty} dh \int_0^{+\infty} dz \frac{d^2 R^{(M)}}{dz dh}(h, z, f) = f. \tag{27}$$

This encodes the fact that when the burst rate is larger than f, individual bursts are not resolved.

The total energy density in gravitational waves produced by cosmic strings will be composed of overlapping signals ($h < h^*$) and nonoverlapping signals, namely bursts ($h > h^*$). The LIGO-Virgo stochastic search pipeline will detect both types of signals. This has been demonstrated for a stochastic background produced by binary neutron stars, whose signals overlap, and binary black holes, whose signals will arrive in a nonoverlapping fashion [61,62]. In this present cosmic string study this effect is negligible: the predicted GW energy density, $\Omega_{\rm GW}^{(M)}$, does not grow significantly (and Fig. 3, top, does not change noticeably) when $h^* \to +\infty$.

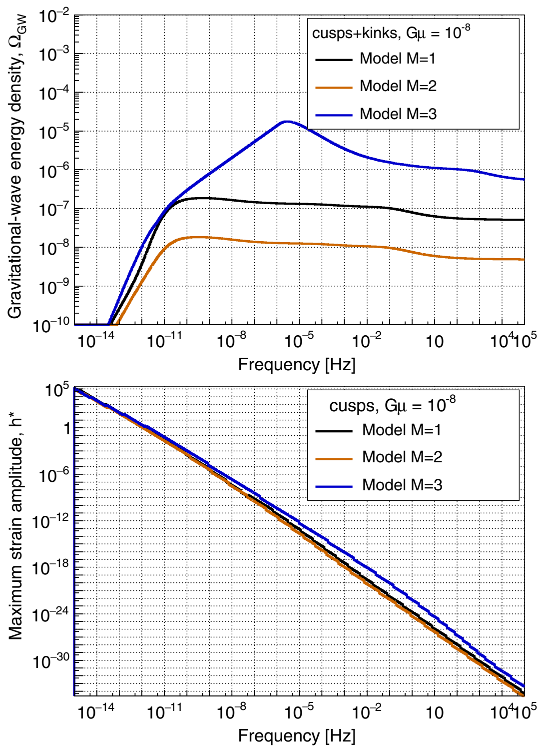

FIG. 3. Top: GW energy density, $\Omega_{\rm GW}^{(M)}(f)$, from cusps and kinks predicted by the three loop distribution models. The string tension $G\mu$ has been fixed to 10^{-8}. Bottom: maximum strain amplitude h^* used for the integration in Eq. (26).

Figure 3 (top) shows the spectra for the three models under consideration, adding both the cusp and the kink contributions and assuming $G\mu = 10^{-8}$. The model 2 spectrum is about 10 times weaker than the spectrum of model 1 over most of the frequency range. As shown in Fig. 4 (top), the spectra are dominated by the contribution of loops in the radiation era over most of the frequency range, including the frequencies accessible to LIGO and Virgo detectors (10–1000 Hz). The difference in normalizations of the loop distributions in the radiation era in the two models, discussed in Sec. II, is therefore the cause for the difference in spectral amplitudes. Note also that at low frequencies ($\sim 10^{-9}$ Hz), at which pulsar timing observations are made, the matter era loops contribute the most.

Figure 3 (top) also shows that the spectrum for model 3 has a significantly higher amplitude than those of models 1 and 2. Figure 4 shows that this spectrum is dominated by the contribution of small loops which, as discussed in Sec. II, are much more numerous in model 3.

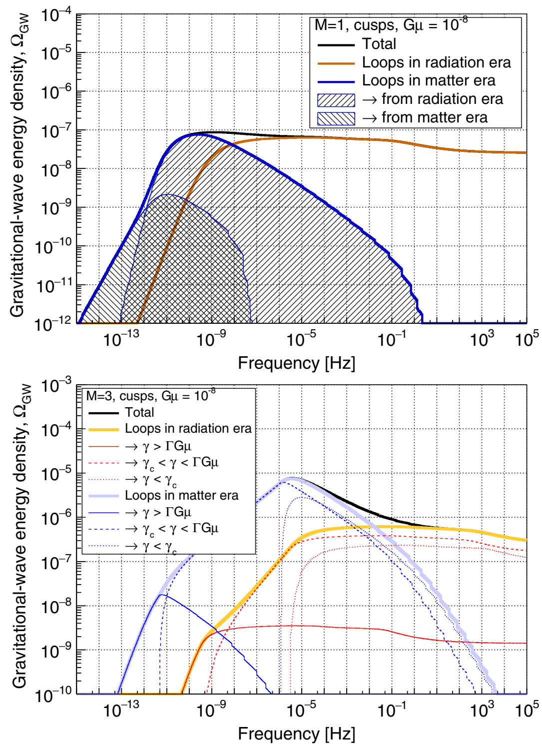

FIG. 4. Top: GW energy density, $\Omega_{\rm GW}^{(M)}(f)$, from cusps for model 1. We have separated the contributions from loops in the radiation ($z > 3366$) and matter ($z < 3366$) eras. Additionally, for loops in the matter era, we have separated the effect of loops produced in the matter era from the ones produced in the radiation era [Eqs. (3), (4), and (6)]. Bottom: GW energy density, $\Omega_{\rm GW}^{(M)}(f)$, from cusps for model 3. The effect of the three loop size regimes is shown [Eqs. (13), (14), and (15)] for the matter and radiation eras.

Figure 3 (bottom) shows the maximum value for the strain amplitude to consider in the integration, h^* as a function of the frequency. At LIGO-Virgo frequencies (10–1000 Hz) the spectrum originates from GWs with strain amplitudes below $\sim 10^{-28}$.

The energy density spectra predicted by the models can be compared with several observational results. First, searches for the stochastic GW background using LIGO and Virgo detectors have been performed, using the initial generation detectors (science run S6, 2009–2010) [63] and the first observation run (O1, 2015–2016) of the advanced detectors [31]. Both searches reported frequency-dependent upper limits on the energy density in GWs. To translate these upper limits into constraints on cosmic string parameters, we define the following likelihood function:

$$\ln L(G\mu, p) \propto \sum_i \frac{-(Y(f_i) - \Omega_{\mathrm{GW}}^{(M)}(f_i; G\mu, p))^2}{\sigma^2(f_i)}, \quad (28)$$

where $Y(f_i)$ and $\sigma(f_i)$ are the measurement and the associated uncertainty of the GW energy density in the frequency bin f_i, and $\Omega_{\mathrm{GW}}^{(M)}(f_i; G\mu, p)$ is the energy density computed by a cosmic string model at the same frequency bin f_i and for some set of model parameters $G\mu$ and p. We evaluate the likelihood function across the parameter space $(G\mu, p)$ and compute the 95% confidence contours for the initial LIGO-Virgo (S6, $41.5 < f < 169$ Hz) [63] and for the most recent Advanced LIGO (O1, $20 < f < 86$ Hz) [31] stochastic background measurements (assuming Bayesian formalism and flat priors in the log parameter space). Since a stochastic background of GWs has not been detected yet, these contours define the excluded regions of the parameter space. We also compute the projected design sensitivity for the Advanced LIGO and Advanced Virgo detectors, using Eq. (28) with $Y(f_i) = 0$ and with the projected $\sigma(f_i)$ for the detector network [64].

Another limit can be computed based on the pulsar timing array (PTA) measurements of the pulse arrival times of millisecond pulsars [29]. This measurement produces a limit on the energy density at nanohertz frequencies—specifically, at 95% confidence $\Omega_{\mathrm{GW}}^{\mathrm{PTA}}(f = 2.8 \times 10^{-9}\ \mathrm{Hz}) < 2.3 \times 10^{-10}$. We directly compare the spectra predicted by our models (at 2.8×10^{-9} Hz) to this constraint.

Finally, indirect limits on the total (integrated over frequency) energy density in GWs can be placed based on the big-bang nucleosynthesis and cosmic microwave background observations. The BBN model and observations of the abundances of the lightest nuclei can be used to constrain the effective number of relativistic degrees of freedom at the time of the BBN, N_{eff}. Under the assumption that only photons and standard light neutrinos contribute to the radiation energy density, N_{eff} is equal to the effective number of neutrinos, corrected for the residual heating of the neutrino fluid due to electron-positron annihilation: $N_{\mathrm{eff}} \simeq 3.046$ [65]. Any deviation from this value can be attributed to extra relativistic radiation, including potentially GWs due to cosmic string kinks and cusps generated prior to BBN. We therefore use the 95% confidence upper limit $N_{\mathrm{eff}} - 3.046 < 1.4$, obtained by comparing the BBN model and the abundances of deuterium and ^{4}He [27], which translates into the following limit on the total energy density in GWs:

$$\Omega_{\mathrm{GW}}^{\mathrm{BBN}}(G\mu, p) = \int_{10^{-10}\ \mathrm{Hz}}^{10^{10}\ \mathrm{Hz}} df\, \Omega_{\mathrm{GW}}^{(M)}(f; G\mu, p) < 1.75 \times 10^{-5}, \quad (29)$$

where the lower bound on the integrated frequency region is determined by the size of the horizon at the time of BBN [60]. In this calculation we only consider kinks and cusps generated before BBN, which implies limiting the redshift integral in Eq. (26) to $z > 5.5 \times 10^9$.

Similarly, presence of GWs at the time of photon decoupling could alter the observed CMB and baryon acoustic oscillation spectra. We apply a similar procedure as in the BBN case, integrating over redshifts before the photon decoupling ($z > 1089$) and over all frequencies above 10^{-15} Hz (horizon size at the time of decoupling) to compute the total energy density of GWs at the time of decoupling. We then compare this quantity to the posterior distribution obtained in [28] to compute the 95% confidence contours:

$$\Omega_{\mathrm{GW}}^{\mathrm{CMB}}(G\mu, p) = \int_{10^{-15}\ \mathrm{Hz}}^{10^{10}\ \mathrm{Hz}} df\, \Omega_{\mathrm{GW}}^{(M)}(f; G\mu, p) < 3.7 \times 10^{-6}, \quad (30)$$

For reference, Fig. 5 shows the energy density spectra for models 1 and 3 using $G\mu = 10^{-8}$. As expected, the

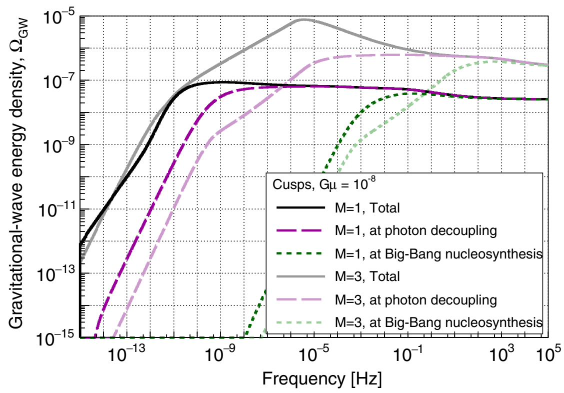

FIG. 5. GW energy density, $\Omega_{\mathrm{GW}}^{(M)}(f)$, from cusps for models 1 and 3. The spectra have been computed at the time of photon decoupling ($z_{\mathrm{CMB}} = 1100$) and at the time of nucleosynthesis ($z_{\mathrm{BBN}} = 5.5 \times 10^9$).

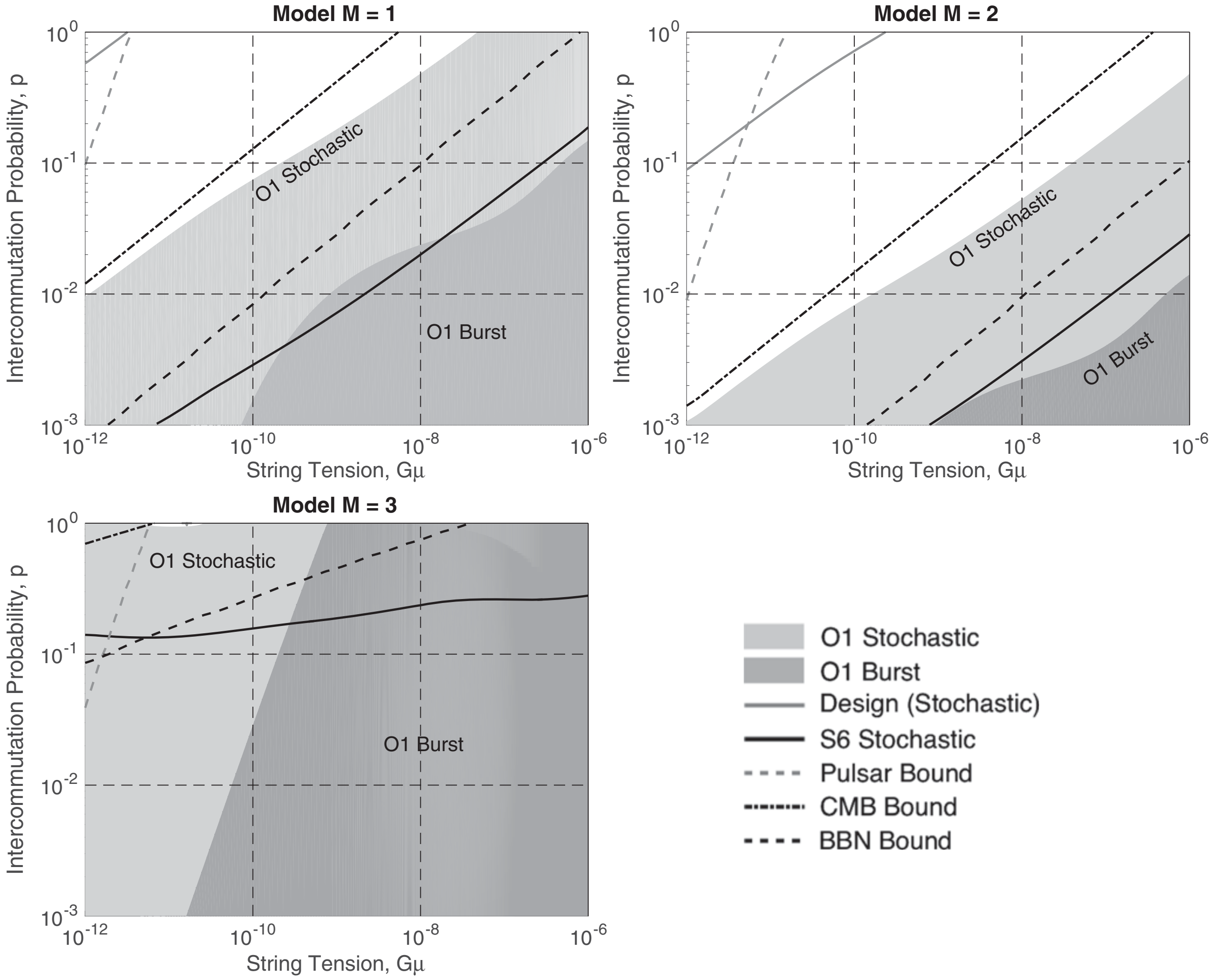

FIG. 6. 95% confidence exclusion regions are shown for three loop distribution models: $M = 1$ (top left), $M = 2$ (top right), and $M = 3$ (bottom left). Shaded regions are excluded by the latest (O1) Advanced LIGO stochastic [31] and burst (presented here) measurements. We also show the bounds from the previous LIGO-Virgo stochastic measurement (S6) [63], from the indirect BBN and CMB bounds [27,28], and from the PTA measurement (pulsar) [29]. Also shown is the projected design sensitivity of the Advanced LIGO and Advanced Virgo experiments (design, stochastic) [64]. The excluded regions are below the respective curves.

contribution from the matter era loops is suppressed at the time of the BBN or of photon decoupling, resulting in the suppression of the spectra at low frequencies. To have negligible systematic errors associated to the numerical integration, we compute Eqs. (29) and (30) using 200 and 250 logarithmically spaced frequency bins respectively.

Figure 6 shows the excluded regions in the parameter spaces of the three models considered here, based on the stochastic observational constraints discussed above.

IV. DISCUSSION

The constraints on the cosmic string tension $G\mu$ and intercommutation probability p are shown in Fig. 6 for the three loop models under consideration: $M = 1$ [8,32] (top left), $M = 2$ [33] (top right), and $M = 3$ [34] (bottom left). We recall that these three models were developed for $p = 1$ and, as explained earlier, for smaller intercommutation probability, we used a $1/p$ dependence for the loop distribution.

The bounds resulting from the burst search performed on O1 data are the least constraining. For model 3 and $p = 1$, the burst search constraint is $G\mu < 8.5 \times 10^{-10}$ at a 95% confidence level. For models 1 and 2, the burst search can only access superstring models ($p < 1$) for which the predicted event rate is larger.

Tighter constraints are obtained when probing the stochastic background of GWs produced by cosmic strings. For model 3, the parameter space studied here is almost entirely excluded by the new constraint derived from the LIGO stochastic O1 analysis. The LIGO stochastic analysis is sensitive to GWs produced in the radiation era. As discussed in Sec. II, in the radiation era, the number of small loops in models 1 and 2 is much smaller than for model 3. When loops are large, the GWs are strongly beamed and the resulting GW detection rate is greatly reduced. As a consequence, experimental bounds using models 1 and 2 are less constraining as can be seen in Fig. 6. For model 1, topological strings ($p = 1$) are constrained by $G\mu < 5 \times 10^{-8}$ with the O1 LIGO stochastic analysis. For model 2, the cosmic string simulation predicts a smaller density of loops and the LIGO constraint is therefore less strict.

In addition to LIGO results, Fig. 6 shows limits from pulsar timing experiments, and indirect limits from BBN and CMB data. These experimental results are complementary as they probe different regions of the loop distributions. The CMB and LIGO stochastic bounds apply for the most part to cosmological loops present in the radiation era ($z > 3300$). The LIGO burst constraint, although weaker, is sensitive to GWs produced in the matter era ($z < 3300$) from loops which themselves were formed in the radiation era. Constraints from pulsar timing experiments are the most competitive. For topological strings, we get $G\mu < 3.8 \times 10^{-12}$, $G\mu < 1.5 \times 10^{-11}$ and $G\mu < 5.7 \times 10^{-12}$ for models 1, 2, and 3 respectively. However, at nanohertz frequencies, they only probe loops formed in the matter era for very small redshifts corresponding to galactic scales ($z \lesssim 10^{-5}$).

The pulsar bound on string parameters will not improve much in the future as the range of strain amplitudes, $10^{-18} \lesssim h \lesssim 10^{-5}$ [see Figs. 3 (bottom) and 8 (right) in Appendix B], allowed by loop models is already fully explored. The indirect bounds from BBN and CMB data will also be limited by the precision on the N_{eff} parameter which can be achieved. The sensitivity of Advanced LIGO detectors, however, will further improve in the coming years. In Fig. 6 we also report the upper limits the stochastic analysis should achieve with an Advanced LIGO-Virgo detector network working at design sensitivity (see also [66,67]). These will probe most of the parameter space for the three models, and, in particular for models 1 and 3, will surpass all of the current bounds.

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APPENDIX A: Λ-CDM COSMOLOGY

In a Λ-CDM universe, the Hubble rate at redshift z is given by

$$H(z) = H_0 \mathcal{H}(z), \tag{A1}$$

where

$$\mathcal{H}(z) = \sqrt{\Omega_\Lambda + \Omega_M (1+z)^3 + \Omega_R \mathcal{G}(z)(1+z)^4}. \tag{A2}$$

We use the latest values of the cosmological parameters [68], $H_0 = 100h\ \mathrm{km\,s^{-1}\,Mpc^{-1}}$, $h = 0.678$, $\Omega_M = 0.308$, $\Omega_R = 9.1476 \times 10^{-5}$, and $\Omega_\Lambda = 1 - \Omega_M - \Omega_R$. At redshift z in the radiation era, the quantity $\mathcal{G}(z)$ is directly related to the effective number of degrees of freedom $g_*(z)$ and the effective number of entropic degrees of freedom $g_S(z)$ by [60]

$$\mathcal{G}(z) = \frac{g_*(z) g_S^{4/3}(0)}{g_*(0) g_S^{4/3}(z)}. \tag{A3}$$

Following [60] we model it by a piecewise constant function whose value changes at the QCD phase transition ($T = 200$ MeV), and at electron-positron annihilation ($T = 200$ keV):

$$\mathcal{G}(z) = \begin{cases} 1 & \text{for } z < 10^9, \\ 0.83 & \text{for } 10^9 < z < 2\times 10^{12}, \\ 0.39 & \text{for } z > 2\times 10^{12}. \end{cases} \tag{A4}$$

Expressions for cosmic time, proper distance, and proper volume element in terms of redshift are given by

$$t(z) = \frac{\varphi_t(z)}{H_0} \quad \text{with} \quad \varphi_t(z) = \int_z^\infty \frac{dz'}{\mathcal{H}(z')(1+z')}, \tag{A5}$$

$$r(z) = \frac{\varphi_r(z)}{H_0} \quad \text{with} \quad \varphi_r(z) = \int_0^z \frac{dz'}{\mathcal{H}(z')}, \tag{A6}$$

$$dV(z) = \frac{\varphi_V(z)}{H_0^3} dz \quad \text{with} \quad \varphi_V(z) = \frac{4\pi\varphi_r^2(z)}{(1+z)^3\mathcal{H}(z)}. \tag{A7}$$

Asymptotically we have

$$\varphi_t(z \ll 1) \sim 0.9566, \tag{A8}$$

$$\varphi_t(z \gg 1) \sim \frac{1}{2\sqrt{\Omega_R \mathcal{G}(z \gg 1)}} z^{-2}, \tag{A9}$$

$$\varphi_r(z \ll 1) \sim z, \tag{A10}$$

$$\varphi_r(z \gg 1) \sim 3.2086. \tag{A11}$$

APPENDIX B: RATE OF GRAVITATIONAL-WAVE BURSTS FROM COSMIC STRINGS

The detection of GWs from cosmic strings is conditioned by the rate of burst events a cosmic string network generates. In this appendix we outline the rate calculation presented in detail in [32] in a form adapted for the three models under consideration.

The expected rate of GW events, observed at frequency f, emitted from a proper volume $dV(z)$ at redshift z, in an interval of amplitudes between A_q and $A_q + dA_q$, and for model M, is given by

$$\frac{d^2 R_q^{(M)}}{dV(z) dA_q}(A_q, z, f) = \frac{1}{1+z} \nu_q^{(M)}(A_q, z) \Delta_q(A_q, z, f), \tag{B1}$$

where Δ_q is the fraction of GW events of amplitude A_q that are observable at frequency f and redshift z. Since cusps emit GW bursts in a cone of solid angle $d\Omega \sim \pi\theta_m^2$ [where θ_m is given in Eq. (20)] and kinks into a fan-shaped set of directions in a solid angle $d\Omega \sim 2\pi\theta_m$, one finds

$$\Delta_q(A_q, z, f) \sim \left(\frac{\theta_m(\ell, z, f)}{2}\right)^{3(2-q)} \times \Theta(1 - \theta_m(z, f, \ell)) \tag{B2}$$

where $\ell = \ell(A_q, z)$ is obtained by inverting Eq. (21).

The number of cusp/kink features per unit spacetime volume on loops with sizes between ℓ and $\ell + d\ell$ is given by

$$\nu_q^{(M)}(\ell, z) d\ell = \frac{2}{\ell} N_q n^{(M)}(\ell, t(z)) d\ell, \tag{B3}$$

where N_q is the number of cusps/kinks per oscillation period. Using Eq. (21) to change variables from ℓ to A_q gives

$$\begin{aligned} \nu_q^{(M)}(A_q, z) dA_q &= \nu_q^{(M)}(\ell(A_q, z), z) \frac{d\ell}{dA_q} dA_q \\ &= \nu_q^{(M)}(\ell(A_q, z), z) \frac{\ell(A_q, z)}{(2-q)A_q} dA_q. \end{aligned} \tag{B4}$$

Injecting the loop distribution of model M, $\mathcal{F}^{(M)}$, Eq. (B1) becomes

$$\begin{aligned} \frac{d^2 R_q^{(M)}}{dz dA_q}(A_q, z, f) = &\frac{2 N_q H_0^{-3} \varphi_V(z)}{(2-q)(1+z) A_q t^4(z)} \\ &\times \mathcal{F}^{(M)}\left(\frac{\ell(A_q, z)}{t(z)}, t(z)\right) \\ &\times \Delta_q(A_q, z, f). \end{aligned} \tag{B5}$$

Alternatively, the rate can also be parametrized by the strain amplitude using Eq. (19):

$$\begin{aligned} \frac{d^2 R_q^{(M)}}{dz dh}(h, z, f) = &\frac{2 N_q H_0^{-3} \varphi_V(z)}{(2-q)(1+z) h t^4(z)} \\ &\times \mathcal{F}^{(M)}\left(\frac{\ell(hf^q, z)}{t(z)}, t(z)\right) \\ &\times \Delta_q(hf^q, z, f). \end{aligned} \tag{B6}$$

The rate of GWs given in Eq. (B6) is marginalized over the strain amplitude and the redshift to compute the GW stochastic background [see Eq. (26)]. The strain amplitude range is limited by two physical conditions: firstly the beaming angle must satisfy $\theta_m < 1$, secondly, in all three models, there is an upper bound for the loop size, $\gamma_{\max}$. These conditions straightforwardly impose [see Eqs. (19) and (21)] that $h_{\min}(z) < h < h_{\max}(z)$, where

$$h_{\min}(z) = \frac{G\mu H_0}{f^2(1+z)\varphi_r(z)}, \tag{B7}$$

$$h_{\max}(z) = \frac{(\gamma_{\max}\varphi_t(z))^{2-q} G\mu}{H_0^{q+1}(1+z)^{q-1} f^q \varphi_r(z)}. \tag{B8}$$

In turn, the condition $h_{\min}(z) \leq h_{\max}(z)$ fixes the upper limit on the redshift, $z_{\max}$. Finally, the overall GW rate is obtained by calculating the double integral:

$$R_q^{(M)} = \int_0^{z_{\max}} dz \int_{h_{\min}(z)}^{h_{\max}(z)} dh \frac{d^2 R_q^{(M)}}{dz dh}(h, z, f). \tag{B9}$$

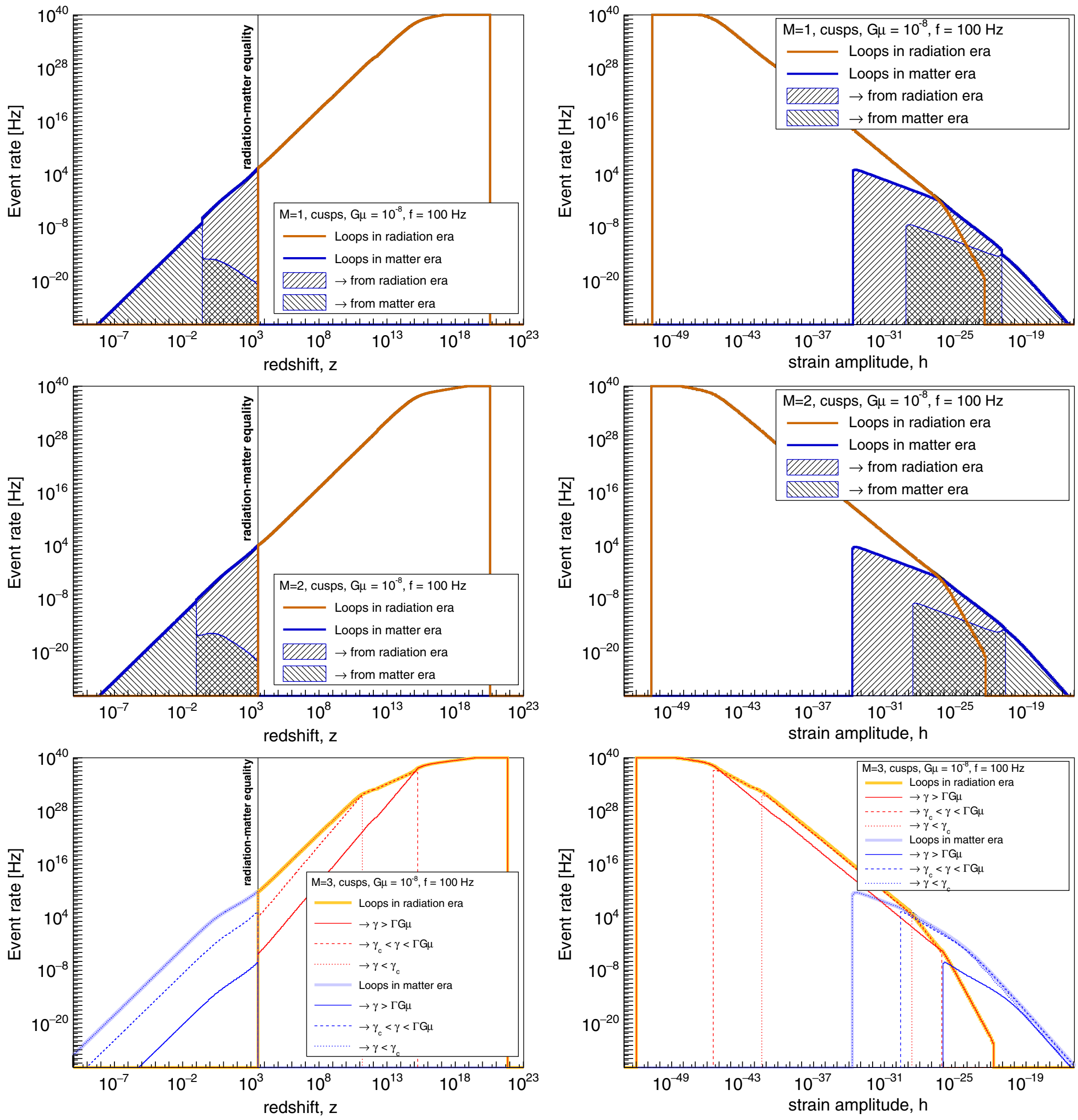

FIG. 7. GW event rate predicted by models $M = 1$ (top row), $M = 2$ (middle row), and $M = 3$ (bottom row) and averaged over either h (left column) or z (right column). The string tension and the wave frequency are fixed to 10^{-8} and to 100 Hz respectively. For models 1 and 2, we separated the contributions from loops in the radiation ($z > 3366$) and matter ($z < 3366$) eras. Additionally, for loops in the matter era, we separated the effect of loops produced in the matter era from the ones produced in the radiation era [Eqs. (3), (4), and (6)]. For model 3, the effect of the three loop size regimes is shown [Eqs. (13), (14), and (15)].

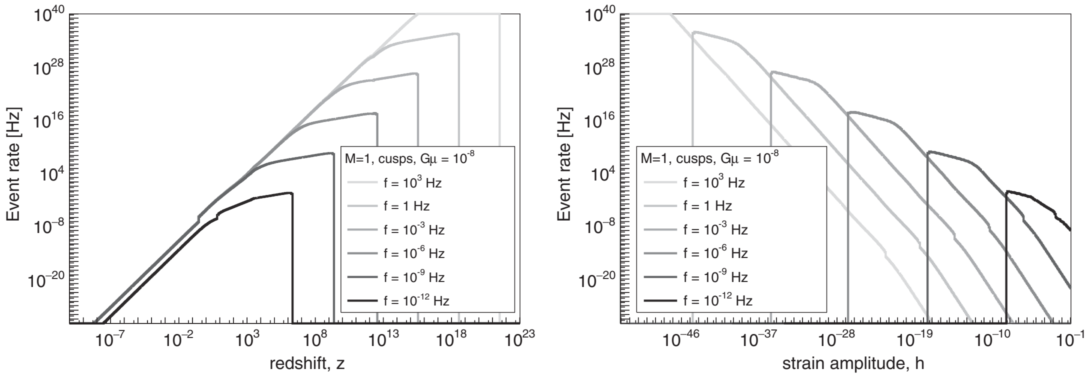

FIG. 8. GW event rate predicted by model $M = 1$ for different frequencies and averaged over either h (left column) or z (right column).

For illustration, we fix $f = 100$ Hz and $G\mu = 10^{-8}$ and we average the GW rate for cusps over either h (left-hand column of Fig. 7) or z (right-hand column of Fig. 7) for models $M = \{1, 2, 3\}$. The contributions from loops in the matter (blue curve) and radiation (red curve) eras are also presented. For loops in the matter era, we separated the effect of loops produced in the radiation era from loops produced in the matter era.

We first observe that all models have the same general dependence on redshift and strain amplitude: high-amplitude GWs are produced in the matter era with a low rate while weak GWs are produced in the radiation era with a high rate. Models differ in the absolute rate of GWs they predict: $R^{(M)}_{\rm cusps}(h = 10^{-23}) = 1.1 \times 10^{-9}$ Hz, 1.2×10^{-10} Hz and 1.0×10^{-6} Hz for models 1, 2, and 3 respectively. As noted in Sec. II, in model 3, small loops are copiously present at all times. Indeed, for model 3, the small loop contribution to the GW rate dominates for $h \gtrsim 10^{-45}$ and $z \lesssim 10^{15}$, while for models 1 and 2 it is negligible.

In Fig. 8, the effect of the wave frequency f is studied ($M = 1$ only). Loops in the radiation era tend to produce high-frequency GWs while low-frequency waves are emitted in the matter era. The event rates presented in Fig. 8 condition the detectability of GWs from cosmic strings using experimental data sets.

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Ng,[87] T. T. Nguyen,[24] D. Nichols,[61] A. B. Nielsen,[9] S. Nissanke,[61,13] A. Nitz,[9] A. Noack,[9] F. Nocera,[28] D. Nolting,[6] M. E. N. Normandin,[97] L. K. Nuttall,[41] J. Oberling,[44] E. Ochsner,[20] E. Oelker,[14] G. H. Ogin,[106] J. J. Oh,[123] S. H. Oh,[123] F. Ohme,[9] M. Oliver,[96] P. Oppermann,[9] Richard J. Oram,[6] B. O'Reilly,[6] R. Ormiston,[42] L. F. Ortega,[5] R. O'Shaughnessy,[115] D. J. Ottaway,[78] H. Overmier,[6] B. J. Owen,[80] A. E. Pace,[82] J. Page,[128] M. A. Page,[60] A. Pai,[109] S. A. Pai,[56] J. R. Palamos,[67] O. Palashov,[120] C. Palomba,[32] A. Pal-Singh,[31] H. Pan,[83] B. Pang,[59] P. T. H. Pang,[87] C. Pankow,[94] F. Pannarale,[95] B. C. Pant,[56] F. Paoletti,[23] A. Paoli,[28] M. A. Papa,[34,20,9] H. R. Paris,[47] W. Parker,[6] D. Pascucci,[43] A. Pasqualetti,[28] R. Passaquieti,[22,23] D. Passuello,[23] B. Patricelli,[135,23] B. L. Pearlstone,[43] M. Pedraza,[1] R. Pedurand,[25,136] L. Pekowsky,[41] A. Pele,[6] S. Penn,[137] C. J. Perez,[44] A. Perreca,[1,102,89] L. M. Perri,[94] H. P. Pfeiffer,[105] M. Phelps,[43] O. J. Piccinni,[90,32] M. Pichot,[62] F. Piergiovanni,[65,66] V. Pierro,[8] G. Pillant,[28] L. Pinard,[25] I. M. Pinto,[8] M. Pitkin,[43] R. Poggiani,[22,23] P. Popolizio,[28] E. K. Porter,[35] A. Post,[9] J. Powell,[43] J. Prasad,[18] J. W. W. Pratt,[33] V. Predoi,[95] T. Prestegard,[20] M. Prijatelj,[9] M. Principe,[8] S. Privitera,[34] R. Prix,[9] G. A. Prodi,[102,89] L. G. Prokhorov,[57] O. Puncken,[9] M. Punturo,[40] P. Puppo,[32] M. Pürrer,[34] H. Qi,[20] J. Qin,[60] S. Qiu,[126] V. Quetschke,[97] E. A. Quintero,[1] R. Quitzow-James,[67] F. J. Raab,[44] D. S. Rabeling,[24] H. Radkins,[44] P. Raffai,[51] S. Raja,[56] C. Rajan,[56] M. Rakhmanov,[97] K. E. Ramirez,[97] P. Rapagnani,[90,32] V. Raymond,[34] M. Razzano,[22,23] J. Read,[27] T. Regimbau,[62] L. Rei,[55] S. Reid,[58] D. H. Reitze,[1,5] H. Rew,[132] S. D. Reyes,[41] F. Ricci,[90,32] P. M. Ricker,[11] S. Rieger,[9] K. Riles,[113] M. Rizzo,[115] N. A. Robertson,[1,43] R. Robie,[43] F. Robinet,[26] A. Rocchi,[17] L. Rolland,[7] J. G. Rollins,[1] V. J. Roma,[67] J. D. Romano,[97] R. Romano,[3,4] C. L. Romel,[44] J. H. Romie,[6] D. Rosińska,[138,52] M. P. Ross,[139] S. Rowan,[43] A. Rüdiger,[9] P. Ruggi,[28] K. Ryan,[44] S. Sachdev,[1] T. Sadecki,[44] L. Sadeghian,[20] M. Sakellariadou,[140] L. Salconi,[28] M. Saleem,[109] F. Salemi,[9] A. Samajdar,[134] L. Sammut,[126] L. M. Sampson,[94] E. J. Sanchez,[1] V. Sandberg,[44] B. Sandeen,[94] J. R. Sanders,[41] B. Sassolas,[25] P. R. Saulson,[41] O. Sauter,[113] R. L. Savage,[44] A. Sawadsky,[21] P. Schale,[67] J. Scheuer,[94] E. Schmidt,[33] J. Schmidt,[9] P. Schmidt,[1,61] R. Schnabel,[31] R. M. S. Schofield,[67] A. Schönbeck,[31] E. Schreiber,[9] D. Schuette,[9,21] B. W. Schulte,[9] B. F. Schutz,[95,9] S. G. Schwalbe,[33] J. Scott,[43] S. M. Scott,[24] E. Seidel,[11] D. Sellers,[6] A. S. Sengupta,[141] D. Sentenac,[28] V. Sequino,[30,17] A. Sergeev,[120] D. A. Shaddock,[24] T. J. Shaffer,[44] A. A. Shah,[128] M. S. Shahriar,[94] L. Shao,[34] B. Shapiro,[47] P. Shawhan,[72] A. Sheperd,[20] D. H. Shoemaker,[14] D. M. Shoemaker,[73] K. Siellez,[73] X. Siemens,[20] M. Sieniawska,[52] D. Sigg,[44] A. D. Silva,[15] A. Singer,[1] L. P. Singer,[76] A. Singh,[34,9,21] R. Singh,[2] A. Singhal,[16,32] A. M. Sintes,[96] B. J. J. Slagmolen,[24] B. Smith,[6] J. R. Smith,[27] R. J. E. Smith,[1] E. J. Son,[123] J. A. Sonnenberg,[20] B. Sorazu,[43] F. Sorrentino,[55] T. Souradeep,[18] A. P. Spencer,[43] A. K. Srivastava,[99] A. Staley,[46] D. A. Steer,[35] M. Steinke,[9] J. Steinlechner,[43,31]

S. Steinlechner,[31] D. Steinmeyer,[9,21] B. C. Stephens,[20] R. Stone,[97] K. A. Strain,[43] G. Stratta,[65,66] S. E. Strigin,[57] R. Sturani,[142] A. L. Stuver,[6] T. Z. Summerscales,[143] L. Sun,[131] S. Sunil,[99] P. J. Sutton,[95] B. L. Swinkels,[28] M. J. Szczepańczyk,[33] M. Tacca,[35] D. Talukder,[67] D. B. Tanner,[5] M. Tápai,[110] A. Taracchini,[34] J. A. Taylor,[128] R. Taylor,[1] T. Theeg,[9] E. G. Thomas,[53] M. Thomas,[6] P. Thomas,[44] K. A. Thorne,[6] K. S. Thorne,[59] E. Thrane,[126] S. Tiwari,[16,89] V. Tiwari,[95] K. V. Tokmakov,[117] K. Toland,[43] M. Tonelli,[22,23] Z. Tornasi,[43] C. I. Torrie,[1] D. Töyrä,[53] F. Travasso,[28,40] G. Traylor,[6] D. Trifirò,[10] J. Trinastic,[5] M. C. Tringali,[102,89] L. Trozzo,[144,23] K. W. Tsang,[13] M. Tse,[14] R. Tso,[1] D. Tuyenbayev,[97] K. Ueno,[20] D. Ugolini,[145] C. S. Unnikrishnan,[111] A. L. Urban,[1] S. A. Usman,[95] H. Vahlbruch,[21] G. Vajente,[1] G. Valdes,[97] M. Vallisneri,[59] N. van Bakel,[13] M. van Beuzekom,[13] J. F. J. van den Brand,[71,13] C. Van Den Broeck,[13] D. C. Vander-Hyde,[41] L. van der Schaaf,[13] J. V. van Heijningen,[13] A. A. van Veggel,[43] M. Vardaro,[49,50] V. Varma,[59] S. Vass,[1] M. Vasúth,[45] A. Vecchio,[53] G. Vedovato,[50] J. Veitch,[53] P. J. Veitch,[78] K. Venkateswara,[139] G. Venugopalan,[1] D. Verkindt,[7] F. Vetrano,[65,66] A. Viceré,[65,66] A. D. Viets,[20] S. Vinciguerra,[53] D. J. Vine,[58] J.-Y. Vinet,[62] S. Vitale,[14] T. Vo,[41] H. Vocca,[39,40] C. Vorvick,[44] D. V. Voss,[5] W. D. Vousden,[53] S. P. Vyatchanin,[57] A. R. Wade,[1] L. E. Wade,[81] M. Wade,[81] R. Walet,[13] M. Walker,[2] L. Wallace,[1] S. Walsh,[20] G. Wang,[16,66] H. Wang,[53] J. Z. Wang,[82] M. Wang,[53] Y.-F. Wang,[87] Y. Wang,[60] R. L. Ward,[24] J. Warner,[44] M. Was,[7] J. Watchi,[91] B. Weaver,[44] L.-W. Wei,[9,21] M. Weinert,[9] A. J. Weinstein,[1] R. Weiss,[14] L. Wen,[60] E. K. Wessel,[11] P. Weßels,[9] T. Westphal,[9] K. Wette,[9] J. T. Whelan,[115] B. F. Whiting,[5] C. Whittle,[126] D. Williams,[43] R. D. Williams,[1] A. R. Williamson,[115] J. L. Willis,[146] B. Willke,[21,9] M. H. Wimmer,[9,21] W. Winkler,[9] C. C. Wipf,[1] H. Wittel,[9,21] G. Woan,[43] J. Woehler,[9] J. Wofford,[115] K. W. K. Wong,[87] J. Worden,[44] J. L. Wright,[43] D. S. Wu,[9] G. Wu,[6] W. Yam,[14] H. Yamamoto,[1] C. C. Yancey,[72] M. J. Yap,[24] Hang Yu,[14] Haocun Yu,[14] M. Yvert,[7] A. Zadrożny,[124] M. Zanolin,[33] T. Zelenova,[28] J.-P. Zendri,[50] M. Zevin,[94] L. Zhang,[1] M. Zhang,[132] T. Zhang,[43] Y.-H. Zhang,[115] C. Zhao,[60] M. Zhou,[94] Z. Zhou,[94] S. J. Zhu,[34,9] X. J. Zhu,[60] M. E. Zucker,[1,14] and J. Zweizig[1]

(LIGO Scientific Collaboration and Virgo Collaboration)

[1]*LIGO, California Institute of Technology, Pasadena, California 91125, USA*
[2]*Louisiana State University, Baton Rouge, Louisiana 70803, USA*
[3]*Università di Salerno, Fisciano, I-84084 Salerno, Italy*
[4]*INFN, Sezione di Napoli, Complesso Universitario di Monte S.Angelo, I-80126 Napoli, Italy*
[5]*University of Florida, Gainesville, Florida 32611, USA*
[6]*LIGO Livingston Observatory, Livingston, Louisiana 70754, USA*
[7]*Laboratoire d'Annecy-le-Vieux de Physique des Particules (LAPP), Université Savoie Mont Blanc, CNRS/IN2P3, F-74941 Annecy, France*
[8]*University of Sannio at Benevento, I-82100 Benevento, Italy and INFN, Sezione di Napoli, I-80100 Napoli, Italy*
[9]*Albert-Einstein-Institut, Max-Planck-Institut für Gravitationsphysik, D-30167 Hannover, Germany*
[10]*The University of Mississippi, University, Mississippi 38677, USA*
[11]*NCSA, University of Illinois at Urbana-Champaign, Urbana, Illinois 61801, USA*
[12]*University of Cambridge, Cambridge CB2 1TN, United Kingdom*
[13]*Nikhef, Science Park, 1098 XG Amsterdam, Netherlands*
[14]*LIGO, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139, USA*
[15]*Instituto Nacional de Pesquisas Espaciais, 12227-010 São José dos Campos, São Paulo, Brazil*
[16]*Gran Sasso Science Institute (GSSI), I-67100 L'Aquila, Italy*
[17]*INFN, Sezione di Roma Tor Vergata, I-00133 Roma, Italy*
[18]*Inter-University Centre for Astronomy and Astrophysics, Pune 411007, India*
[19]*International Centre for Theoretical Sciences, Tata Institute of Fundamental Research, Bengaluru 560089, India*
[20]*University of Wisconsin-Milwaukee, Milwaukee, Wisconsin 53201, USA*
[21]*Leibniz Universität Hannover, D-30167 Hannover, Germany*
[22]*Università di Pisa, I-56127 Pisa, Italy*
[23]*INFN, Sezione di Pisa, I-56127 Pisa, Italy*
[24]*OzGrav, Australian National University, Canberra, Australian Capital Territory 0200, Australia*
[25]*Laboratoire des Matériaux Avancés (LMA), CNRS/IN2P3, F-69622 Villeurbanne, France*
[26]*LAL, Univ. Paris-Sud, CNRS/IN2P3, Université Paris-Saclay, F-91898 Orsay, France*
[27]*California State University Fullerton, Fullerton, California 92831, USA*
[28]*European Gravitational Observatory (EGO), I-56021 Cascina, Pisa, Italy*
[29]*Chennai Mathematical Institute, Chennai 603103, India*

[30]Università di Roma Tor Vergata, I-00133 Roma, Italy
[31]Universität Hamburg, D-22761 Hamburg, Germany
[32]INFN, Sezione di Roma, I-00185 Roma, Italy
[33]Embry-Riddle Aeronautical University, Prescott, Arizona 86301, USA
[34]Albert-Einstein-Institut, Max-Planck-Institut für Gravitationsphysik, D-14476 Potsdam-Golm, Germany
[35]APC, AstroParticule et Cosmologie, Université Paris Diderot, CNRS/IN2P3, CEA/Irfu, Observatoire de Paris, Sorbonne Paris Cité, F-75205 Paris Cedex 13, France
[36]Korea Institute of Science and Technology Information, Daejeon 34141, Korea
[37]West Virginia University, Morgantown, West Virginia 26506, USA
[38]Center for Gravitational Waves and Cosmology, West Virginia University, Morgantown, West Virginia 26505, USA
[39]Università di Perugia, I-06123 Perugia, Italy
[40]INFN, Sezione di Perugia, I-06123 Perugia, Italy
[41]Syracuse University, Syracuse, New York 13244, USA
[42]University of Minnesota, Minneapolis, Minnesota 55455, USA
[43]SUPA, University of Glasgow, Glasgow G12 8QQ, United Kingdom
[44]LIGO Hanford Observatory, Richland, Washington 99352, USA
[45]Wigner RCP, RMKI, H-1121 Budapest, Konkoly Thege Miklós út 29-33, Hungary
[46]Columbia University, New York, New York 10027, USA
[47]Stanford University, Stanford, California 94305, USA
[48]Università di Camerino, Dipartimento di Fisica, I-62032 Camerino, Italy
[49]Università di Padova, Dipartimento di Fisica e Astronomia, I-35131 Padova, Italy
[50]INFN, Sezione di Padova, I-35131 Padova, Italy
[51]MTA Eötvös University, "Lendulet" Astrophysics Research Group, Budapest 1117, Hungary
[52]Nicolaus Copernicus Astronomical Center, Polish Academy of Sciences, 00-716, Warsaw, Poland
[53]University of Birmingham, Birmingham B15 2TT, United Kingdom
[54]Università degli Studi di Genova, I-16146 Genova, Italy
[55]INFN, Sezione di Genova, I-16146 Genova, Italy
[56]RRCAT, Indore MP 452013, India
[57]Faculty of Physics, Lomonosov Moscow State University, Moscow 119991, Russia
[58]SUPA, University of the West of Scotland, Paisley PA1 2BE, United Kingdom
[59]Caltech CaRT, Pasadena, California 91125, USA
[60]OzGrav, University of Western Australia, Crawley, Western Australia 6009, Australia
[61]Department of Astrophysics/IMAPP, Radboud University Nijmegen, P.O. Box 9010, 6500 GL Nijmegen, Netherlands
[62]Artemis, Université Côte d'Azur, Observatoire Côte d'Azur, CNRS, CS 34229, F-06304 Nice Cedex 4, France
[63]Institut de Physique de Rennes, CNRS, Université de Rennes 1, F-35042 Rennes, France
[64]Washington State University, Pullman, Washington 99164, USA
[65]Università degli Studi di Urbino 'Carlo Bo', I-61029 Urbino, Italy
[66]INFN, Sezione di Firenze, I-50019 Sesto Fiorentino, Firenze, Italy
[67]University of Oregon, Eugene, Oregon 97403, USA
[68]Laboratoire Kastler Brossel, UPMC-Sorbonne Universités, CNRS, ENS-PSL Research University, Collège de France, F-75005 Paris, France
[69]Carleton College, Northfield, Minnesota 55057, USA
[70]Astronomical Observatory Warsaw University, 00-478 Warsaw, Poland
[71]VU University Amsterdam, 1081 HV Amsterdam, Netherlands
[72]University of Maryland, College Park, Maryland 20742, USA
[73]Center for Relativistic Astrophysics and School of Physics, Georgia Institute of Technology, Atlanta, Georgia 30332, USA
[74]Université Claude Bernard Lyon 1, F-69622 Villeurbanne, France
[75]Università di Napoli 'Federico II', Complesso Universitario di Monte S.Angelo, I-80126 Napoli, Italy
[76]NASA Goddard Space Flight Center, Greenbelt, Maryland 20771, USA
[77]RESCEU, University of Tokyo, Tokyo, 113-0033, Japan
[78]OzGrav, University of Adelaide, Adelaide, South Australia 5005, Australia
[79]Tsinghua University, Beijing 100084, China
[80]Texas Tech University, Lubbock, Texas 79409, USA
[81]Kenyon College, Gambier, Ohio 43022, USA
[82]The Pennsylvania State University, University Park, Pennsylvania 16802, USA
[83]National Tsing Hua University, Hsinchu City, 30013 Taiwan, Republic of China

[84]*Charles Sturt University, Wagga Wagga, New South Wales 2678, Australia*
[85]*University of Chicago, Chicago, Illinois 60637, USA*
[86]*Pusan National University, Busan 46241, Korea*
[87]*The Chinese University of Hong Kong, Shatin, NT, Hong Kong*
[88]*INAF, Osservatorio Astronomico di Padova, Vicolo dell'Osservatorio 5, I-35122 Padova, Italy*
[89]*INFN, Trento Institute for Fundamental Physics and Applications, I-38123 Povo, Trento, Italy*
[90]*Università di Roma 'La Sapienza', I-00185 Roma, Italy*
[91]*Université Libre de Bruxelles, Brussels 1050, Belgium*
[92]*Sonoma State University, Rohnert Park, California 94928, USA*
[93]*Montana State University, Bozeman, Montana 59717, USA*
[94]*Center for Interdisciplinary Exploration & Research in Astrophysics (CIERA), Northwestern University, Evanston, Illinois 60208, USA*
[95]*Cardiff University, Cardiff CF24 3AA, United Kingdom*
[96]*Universitat de les Illes Balears, IAC3—IEEC, E-07122 Palma de Mallorca, Spain*
[97]*The University of Texas Rio Grande Valley, Brownsville, Texas 78520, USA*
[98]*Bellevue College, Bellevue, Washington 98007, USA*
[99]*Institute for Plasma Research, Bhat, Gandhinagar 382428, India*
[100]*The University of Sheffield, Sheffield S10 2TN, United Kingdom*
[101]*California State University, Los Angeles, 5151 State University Dr, Los Angeles, California 90032, USA*
[102]*Università di Trento, Dipartimento di Fisica, I-38123 Povo, Trento, Italy*
[103]*Montclair State University, Montclair, New Jersey 07043, USA*
[104]*National Astronomical Observatory of Japan, 2-21-1 Osawa, Mitaka, Tokyo 181-8588, Japan*
[105]*Canadian Institute for Theoretical Astrophysics, University of Toronto, Toronto, Ontario M5S 3H8, Canada*
[106]*Whitman College, 345 Boyer Avenue, Walla Walla, Washington 99362 USA*
[107]*School of Mathematics, University of Edinburgh, Edinburgh EH9 3FD, United Kingdom*
[108]*University and Institute of Advanced Research, Gandhinagar Gujarat 382007, India*
[109]*IISER-TVM, CET Campus, Trivandrum Kerala 695016, India*
[110]*University of Szeged, Dóm tér 9, Szeged 6720, Hungary*
[111]*Tata Institute of Fundamental Research, Mumbai 400005, India*
[112]*INAF, Osservatorio Astronomico di Capodimonte, I-80131, Napoli, Italy*
[113]*University of Michigan, Ann Arbor, Michigan 48109, USA*
[114]*American University, Washington, DC 20016, USA*
[115]*Rochester Institute of Technology, Rochester, New York 14623, USA*
[116]*University of Białystok, 15-424 Białystok, Poland*
[117]*SUPA, University of Strathclyde, Glasgow G1 1XQ, United Kingdom*
[118]*University of Southampton, Southampton SO17 1BJ, United Kingdom*
[119]*University of Washington Bothell, 18115 Campus Way NE, Bothell, Washington 98011, USA*
[120]*Institute of Applied Physics, Nizhny Novgorod, 603950, Russia*
[121]*Seoul National University, Seoul 08826, Korea*
[122]*Inje University Gimhae, South Gyeongsang 50834, Korea*
[123]*National Institute for Mathematical Sciences, Daejeon 34047, Korea*
[124]*NCBJ, 05-400 Świerk-Otwock, Poland*
[125]*Institute of Mathematics, Polish Academy of Sciences, 00656 Warsaw, Poland*
[126]*OzGrav, School of Physics and Astronomy, Monash University, Clayton 3800, Victoria, Australia*
[127]*Hanyang University, Seoul 04763, Korea*
[128]*NASA Marshall Space Flight Center, Huntsville, Alabama 35811, USA*
[129]*ESPCI, CNRS, F-75005 Paris, France*
[130]*Southern University and A&M College, Baton Rouge, Louisiana 70813, USA*
[131]*OzGrav, University of Melbourne, Parkville, Victoria 3010, Australia*
[132]*College of William and Mary, Williamsburg, Virginia 23187, USA*
[133]*Indian Institute of Technology Madras, Chennai 600036, India*
[134]*IISER-Kolkata, Mohanpur, West Bengal 741252, India*
[135]*Scuola Normale Superiore, Piazza dei Cavalieri 7, I-56126 Pisa, Italy*
[136]*Université de Lyon, F-69361 Lyon, France*
[137]*Hobart and William Smith Colleges, Geneva, New York 14456, USA*
[138]*Janusz Gil Institute of Astronomy, University of Zielona Góra, 65-265 Zielona Góra, Poland*
[139]*University of Washington, Seattle, Washington 98195, USA*
[140]*King's College London, University of London, London WC2R 2LS, United Kingdom*
[141]*Indian Institute of Technology, Gandhinagar Ahmedabad Gujarat 382424, India*

[142]*International Institute of Physics, Universidade Federal do Rio Grande do Norte, Natal RN 59078-970, Brazil*
[143]*Andrews University, Berrien Springs, Michigan 49104, USA*
[144]*Università di Siena, I-53100 Siena, Italy*
[145]*Trinity University, San Antonio, Texas 78212, USA*
[146]*Abilene Christian University, Abilene, Texas 79699, USA*

[†]Deceased.