Hidden-Sector Spectroscopy with Gravitational Waves from Binary Neutron Stars

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ABSTRACT

We show that neutron star binaries can be ideal laboratories to probe hidden sectors with a long range force. In particular, it is possible for gravitational wave detectors such as LIGO and Virgo to resolve the correction of waveforms from ultralight dark gauge bosons coupled to neutron stars. We observe that the interaction of the hidden sector affects both the gravitational wave frequency and amplitude in a way that cannot be fitted by pure gravity.

1. INTRODUCTION

The LIGO and Virgo collaborations have observed gravitational waves (GWs) GW170817 from an inspiraling binary of neutron stars (NSs) Abbott et al. (2017a). This signal, and the associated electromagnetic counterpart GRB170817a Abbott et al. (2017b,c), provide a wealth of information about neutron stars. In particular, the finely measured waveform of GWs with relatively long duration (∼20 min) could reveal many otherwise hidden information about the physical properties of NSs themselves and of their ambient environment Randall & Xianyu (2018), many of which could be originated from a new physics sector Ellis et al. (2017).

In this Letter, we show that such binary NSs can be an ideal laboratory for probing a long range force mediated by a new ultralight particle with inverse mass m−1 comparable to or larger than the binary separation r in natural units. The LIGO band for binary separation r is roughly O(10–1000) km, and this translates to O(10−13 − 10−11) eV for mass. A prototypical example of this long range force is a hidden sector with asymmetric dark matter (aDM) along with an ultralight mediator, e.g. a dark photon. The charged aDM particles can be trapped by a NS during its lifetime or could have been present in the star at birth. Binary NSs containing such aDM then feel a long range force mediated by the massive but ultralight dark photon.

Independent of the detailed mechanism of aDM-trapping in NSs, we shall show, in a fairly model-independent manner, that clean and detectable GW signals can be generated by aDM-carrying neutron star (NS) pairs. In the current work, we only make two mild assumptions about the hidden-sector companion of NSs: 1) For the mechanism to work, it is assumed that the DM comprises a small mass fraction of the NSs. This implies that the attractive gravitational force is stronger than the hidden repulsion on relevant macroscopic scales. 2) The mass profile of DM in NSs does not differ significantly from the neutron star matter, as is the case for a hidden sector without a repulsive interaction Rezaei (2017). This ensures that the point mass approximation for binary NSs holds, and that clean and well-defined GW signals can be derived.

The chirping GWs generated by a pair of purely self-gravitating point masses m1 and m2 has the following time-dependent frequency fGW, to leading order in the post-Newtonian expansion,

\[ f_{\text{GW}}(t) = \frac{1}{\pi} \left( \frac{G m_{c}}{c^3} \right)^{-5/8} \left( \frac{5}{256} \frac{1}{t_0 - t} \right)^{3/8}, \]

where \( t_0 \) is the time of the coalescence and \( m_{c} = (m_1 m_2)^{3/5}/(m_1 + m_2)^{1/5} \) is the chirp mass. For GW170817, the chirp has been measured with rather high precision as \( m_{c} = 1.188^{+0.004}_{-0.002} M_{\odot} \) Abbott et al. (2017a).

The chirp signal is subject to corrections from hidden sectors, of which we identify two important effects with the above mild phenomenological assumptions. The first effect is due to the Yukawa potential between the two NSs coming from the exchange of a massive dark photon. In the case that dark photon mass \( m_{\gamma} \) lies within the LIGO band, the Yukawa repulsion between to two NSs is virtually absent when the binary separation \( r \gg m_{\gamma}^{-1} \), but behaves like Coulomb repulsion when \( r \ll m_{\gamma}^{-1} \). The Coulomb repulsion would affect the observed chirp mass of the binary, and therefore the Yukawa potential will generate a characteristic shift in chirp mass during...
the inspiraling phase that can in principle be observed by LIGO and Virgo detectors.

The second effect is from the fact that such inspiraling NS binaries generate dark radiation, so long as the wavelength of the radiation is much greater than $m_{\gamma}^{-1}$. If the two NSs do not share exactly the same (dark) charge-mass ratio, then the dark radiation develops an electric dipole component which is qualitatively different from GWs because the latter start only from quadrupole level, and therefore the dark dipole radiation will generate a distinct correction to GW signals.

Conversely, in detected events without observed hidden-sector corrections such as in GW170817, bounds can be placed on the relative strength of the long range force between the two NSs, as will be detailed below. Importantly, these constraints hold independently of both the strength and form of the portal coupling, and the DM relic density.

In the rest of this Letter we first describe a generic model of asymmetric dark matter with an ultralight mediator. We then describe in detail the two effects, the dark repulsion and the dark radiation. We summarize with a discussion on the applicability of our results.

2. A GENERIC MODEL

Though our analysis of hidden sector corrections to GWs is rather generic and model independent, it is very helpful to illustrate the general point with a simple model of DM charged under $U(1)'$,

$$\mathcal{L} = \mathcal{L}_V + \mathcal{L}_\chi + \mathcal{L}_{mix},$$

where $\mathcal{L}_V$ is the vector potential of the gauge fields, and $\mathcal{L}_\chi$ the Lagrangian of DM,

$$\mathcal{L}_V = -\frac{1}{4} V_{\mu\nu} V^{\mu\nu} + \frac{1}{2} m_{\gamma}^2 V_\mu V^\mu$$

$$\mathcal{L}_\chi = \bar{\chi}(\gamma^\mu(i\partial_\mu - g'V_\mu) - m_\chi) \chi.$$  \hspace{1cm} (4)

Here $m_\gamma$ is the dark photon Stuckelberg mass term, $m_\chi$ is the dark matter mass, $g'$ is the gauge coupling. We will exchange the latter for $\alpha' = g'^2/4\pi$ throughout the letter. We make no assumptions on the form of $\mathcal{L}_{mix}$, though traditionally that part of the lagrangian holds the photon-dark photon kinetic mixing terms. Throughout, we require the dark photon mass to reside within or be lighter than the LIGO-Virgo detection window, $m_{\gamma}^{-1} > \mathcal{O}(10)$ km.

3. DARK REPULSION

Due to the presence of aDM, the two NSs carry like $U(1)'$ charges, such that a repulsive force is generated. The range of the force is determined by the mass of the dark photon. When the distance of the two NSs is large the interaction can be neglected; when they are within the effective range $m_{\gamma}^{-1}$, the repulsive dark force is switched on, and effectively behaves like a Coulomb force at small distance. This is the reason NS mergers are particularly good for constraining dark photons in the range $m_{\gamma}^{-1} \sim \mathcal{O}(10 - 1000)$ km, with NS radius being $\mathcal{O}(10)$ km, and separation when the signal enters LIGO-Virgo window up to $\mathcal{O}(1000)$ km.

For a binary system of NSs with mass $m_{1,2}$ and dark $U(1)'$ charge $q_{1,2}$, the orbital frequency of the inspiraling binary is given by

$$\omega^2 = \frac{Gm}{r^3} \left( 1 - \alpha'e^{-m_{\gamma}r} \right),$$

where $m = m_1 + m_2$ is the total mass of the binary, $r$ is the binary separation, and $\alpha'$ is defined to be,

$$\alpha' \equiv \frac{\alpha' q_1 q_2}{G m_1 m_2} = \frac{\alpha'}{G} \left( \frac{f}{m_\chi} \right)^2,$$

where $f$ is the fraction of aDM in each NS. The total energy of the system is,

$$E = E_G + E_V = -\frac{Gm\mu}{2r} \left( 1 - \alpha'e^{-m_{\gamma}r} \right),$$

where $\mu = m_1 m_2 / m$ is the reduced mass.

The result of the new force gives a distortion of the waveform due to the fact that the $r$ dependence of Yukawa potential is different from gravity. To see the point qualitatively, we first note that the Yukawa repulsion is absent when $r > m_{\gamma}^{-1}$, which corresponds to the early stage of inspiraling in the LIGO band. In this regime the waveform is identical to a purely self-gravitating pair as shown in (1). On the other hand, when the binary separation reduces below $m_{\gamma}^{-1}$, which corresponds to a later stage in LIGO band, the dark repulsion can be approximated by a Coulomb force with $r^{-2}$ law. This force is identical with gravity in $r$ dependence and thus the waveform will still be (1) but with a modified chirp mass due to the Coulomb repulsion. In summary, the apparent chirp mass is,

$$\hat{m}_c = \begin{cases} m_c & (r > m_{\gamma}^{-1}), \\ \left( 1 - \alpha' \right)^{2/5} m_c & (r < m_{\gamma}^{-1}). \end{cases}$$

The factor $(1 - \alpha')^{2/5}$ in the second line is found by solving the equation of energy conservation $E = -P_{GW}$, where $E$ is given in (7) with $m_{\gamma} \to 0$ and $P_{GW}$ is the power of quadrupole radiation of GWs and is given by (see, e.g., Maggiore (2007) for pedagogical introductions and Randall & Xianyu (2018) for a short review),

$$P_{GW} = \frac{32G\mu^2 \omega^4 r^4}{5c^5}.$$
Therefore, we see that the net effect of Yukawa repulsion is to generate a characteristic and observable shift in the chirp mass, if the dark photon mass lies within the LIGO band. We stress that the shift of the chirp mass is not a uniform or time-independent correction, but occurs within a short period of time in the LIGO band when $r$ goes across the Yukawa threshold and thus is resolvable by a precise measurement of the chirp mass.

The full waveform corrected by Yukawa repulsion can be worked out again by solving the energy-conservation equation $\dot{E} = -P_{GW}$ but with the full expression of $E$ in (7). Here it is more convenient to work out $r(t)$ instead of $\omega(t)$. This can be solved analytically to first order in $\hat{\alpha}'$,

$$\frac{d\hat{r}}{dt} \simeq \frac{64G^3m^2\mu}{5e^3r^5} \left[ 1 + \hat{\alpha}' e^{-m_\nu r}(m_\nu r - 2) \right].$$

From this equation we can solve for $r(t)$ to first order in $\hat{\alpha}'$. Let $r(0) = r_0$ be the binary separation when the signal enters the LIGO-Virgo band then $r(t)$ can be found

$$f(r) \equiv \frac{e^{-m_\nu r}}{m_\nu^4} \left[ 12 + m_\nu r(2 + m_\nu r)(6 + m_\nu^2r^2) \right].$$

With $r(t)$ known, the corrected time dependence of the orbital frequency $\omega(t)$ can be found from (5), and the corrected chirp signal has frequency $f_{GW}$ and amplitude $A_{GW}$ Maggiore (2007),

$$f_{GW}(t) = \frac{\omega(t)}{\pi},$$

$$A_{GW}(t) = \frac{12G}{d_L c^2} \cdot 2\mu \omega^2(t)^2(t),$$

where $d_L$ is the luminosity distance of the source. The effect on frequency $f_{GW}(t)$ and amplitude $A_{GW}(t)$ of the GW signal are shown as solid curves in Figure 1 and Figure 2 respectively. The initial frequency is set to be consistent with Ref. Abbott et al. (2017a), which corresponds to different separations of the two stars. Although the most significant effect of dark repulsion in Figure 1 is the strengthened signal duration, we stress again that it is the less visible effect of the shifted chirp mass that lead to a distinct signal non-degenerate with the pure gravity case (black curve). Furthermore, we emphasize that the value of $\hat{\alpha}'$ taken in Figure 1 and Figure 2 and much exaggerated only for illustrative purpose. In reality, the value of $\hat{\alpha}'$ is much smaller and we shall show later that the LIGO-Virgo measurement of GW170817 can actually be sensitive to a smaller $\hat{\alpha}'$ down to $\mathcal{O}(10^{-2})$.

![Figure 1.](image-url)  

**Figure 1.** The change of frequency-time dependence due to the dark boson. Here, the dimensionless coupling in (6) is set to $\hat{\alpha}' = 0.5$ for illustration. (See main text for a realistic choice of $\hat{\alpha}'$ and the description of the signal in the change of shape.) The NS masses are chosen to be $m_1 = 1.51 M_\odot$ and $m_2 = 1.24 M_\odot$ respectively, such that the chirp mass before aDM correction corresponds to that of GW170817. The corresponding limits $m_\nu \rightarrow \infty$ and $\hat{\alpha}' \rightarrow 0$ are degenerate and given by the black curve. The solid lines correspond to zero charge-mass difference $\gamma$ and thus show the effect of dark repulsion. The dotted lines correspond to a charge-mass difference of $\gamma = q_1$, and therefore show the competing effects of dark dipole radiation and repulsion.

It is seen that the correction due to dark dipole radiation is important for relatively light dark mediators, and is switched on for $m_\nu c/\omega < 1$ as detailed in the text.

4. DARK RADIATION

A separate effect of the hidden sector comes from the fact that a pair of dark charged inspiraling NSs radiate dark photons. If the charge-mass ratios of the two stars are not identical, a net (dark electric) dipole moment stimulates dipole radiation. This dark dipole radiation drains additional energy away from the system and thus affects the waveform of the chirp signal. Importantly, the dipole radiation is stimulated only when the frequency is higher than the dark photon mass, i.e. $m_\nu c/\omega < 1$, and is otherwise quenched. Therefore, the dark photon mass should be $m_\nu^{-1} > \mathcal{O}(10^2)$km for the dipole radiation to be generated within the LIGO band. This includes a very light dark photon with $m_\nu^{-1} > \mathcal{O}(1000\text{km})$ where the dark photon becomes effectively massless in the Yukawa potential for the whole range of LIGO band, and the dipole radiation becomes the most important correction. In this case the dark repulsion is well approximated by a Coulomb potential, such that the repulsive effect is degenerate with the pure gravity case with modified chirp mass $\tilde{m}_c = (1 - \hat{\alpha}')^{2/5} m_c$ and therefore cannot be observed directly.
The total power of dark dipole radiation is given by Krause et al. (1994),
\[
P_{\text{dark}} = \begin{cases} 
\frac{\alpha' \gamma^2}{3c^3} \sqrt{1 - \left(\frac{m_V c}{\omega}\right)^2} & (\omega > m_V c), \\
0 & (\omega \leq m_V c).
\end{cases}
\]
where \(\gamma \equiv \mu |\rho_1 - \rho_2|\) and where \(\rho_i = g_i/m_i\) \((i = 1, 2)\) is the charge-mass ratio of each NS. From this it is seen that the effect is absent if both NS carry the same DM mass fraction, as expected. The chirp signal is now found from \(\dot{\omega} = -P_{GW} - P_{\text{dark}},\) or more explicitly,
\[
\frac{d\omega}{dt} = X \omega^{11/3} + Y \omega^3,
\]
where
\[
X = \frac{96}{5} \left(\frac{\tilde{G} \hat{m}_v}{c^3}\right)^{5/3}, \quad Y = \frac{\alpha' \gamma^2}{c^3 \mu}.
\]
The equation can be readily integrated to get the chirp signal \(\omega(t)\) and it is clear that the resulted waveform deviates from the standard form in equation (1). In particular, if we assume the correction is small (which must be for detected events), then we are allowed to expand the result in terms of small \(\gamma\) to get the following modified chirp signal,
\[
\omega = \left(\frac{3}{8X(t_c - t)}\right)^{3/8} - \frac{Y}{10X} \left(\frac{3}{8X(t_c - t)}\right)^{1/8},
\]
where \(t_c\) is the time of coalescence. With \(\omega(t)\) and \(r(t)\) known, the frequency and amplitude of the signal can again be obtained from (13) and (14), which is shown as dotted curves in Figure 1 and Figure 2.

When the leading electric dipole is suppressed by the small difference in the charge-mass-ratio, the higher multipoles start to dominate, including the magnetic dipole and electric quadrupole. Again, they are stimulated only if \(m_V c/\omega < 1\). In this case, the total power of the magnetic dipole radiation is,
\[
P_{M1} = \frac{\alpha' \mu^2}{12c^5 m_i^2} (m_2 \rho_1 + m_1 \rho_2)^2 \omega^6 r^4 \times \left[1 - \left(\frac{m_V c}{\omega}\right)^2\right]^{3/2}.
\]
Whereas the total power of the electric quadrupole radiation is,
\[
P_{E2} = \frac{\alpha' \mu^2}{72c^5 m_i^2} (m_2 \rho_1 + m_1 \rho_2)^2 \omega^6 r^4 \times \left[1 - \left(\frac{m_V c}{\omega}\right)^2\right]^{3/2} \left[1 + \frac{22}{15} \left(\frac{m_V c}{\omega}\right)^2\right]^{3/2}.
\]

Both are zero when \(\omega \leq m_V c\). Not surprisingly, they are degenerate with the leading order GW radiation and thus do not generate independent corrections to the original chirp signal.

5. COMPARISON WITH SCALAR TENSOR THEORY

It is known that a large class of gravitational interactions that deviate from the general relativity prediction can be described in terms of the scalar-tensor theory. As one type of scalar-tensor theory, Brans-Dicke theory has constant scalar charge. Therefore, it receives universal bounds and can potentially affect the NS’s equation of state Will & Zaglauer (1989). The current bounds on Brans-Dicke theory come from solar system tests and lack of dipole radiation from certain binary pulsar systems. It is believed that within in current bound Brans-Dicke theory is not likely to generate a dipole radiation large enough to be observed in GW experiments Sampson et al. (2014). However, there is another class of scalar-tensor theories in which the scalar charge can be dynamically acquired Damour & Esposito-Farese (1992, 1993); Barausse et al. (2013); Damour & Esposito-Farese (1996); Palenzuela et al. (2014). In particular, Ref. Barausse et al. (2013) shows a way that the scalar charge can be dynamically generated in the binary merger system after the orbital binding energy reaches a threshold. Rigorous bounds on the theory parameter \(\beta_{ST}\) from NS mergers are demonstrated in Ref. Sampson et al. (2014) for massless scalar case and Ref. Sagunski et al. (2018) for massive scalar. In particular, the scalar charge defined in Ref. Sagunski et al. (2018) can be mapped to our scalar charge as \(\alpha' = -\alpha^2 = -2\beta^2\). For example, in \(f(R)\) gravity, \(\beta = 6^{-1/2}\) corresponds to \(\alpha' = -1/3\). While Sagunski et al. (2018) considers the scalar interac-
tation being attractive, we consider the repulsive case due to NS binary carrying like dark charge, i.e. $\tilde{\alpha}' > 0$. In addition, even though Ref. Sampson et al. (2014) studies a massless scalar, the acquired scalar change in the inspiral phase effectively switches on the interaction, which is similar to the effect in our model. We note that, while in the scalar tensor theory one can have both dipole radiation and ‘turn-on’ signal in the inspiral phase due to scalarization, which can be captured with ppEθ template Sampson et al. (2014, 2013), our model only permits either of the two signals at a time, depending on the mass of the mediator. We also note in passing that, another difference between the scalar tensor theory and our model lies in the fact that the dark force in our model would not affect the black hole mergers.

6. DISCUSSIONS

We can distinguish two different observational windows for our binary neutron star “spectrometer” depending on the mediator mass. Firstly, in the range of $\mathcal{O}(10)$ km $\lesssim (m_V)^{-1} \lesssim \mathcal{O}(1000)$ km, for which dark radiation is mostly suppressed. As can be seen from (7), the dark repulsion is suppressed in the early stages of the merger, but behaves effectively as Coulomb repulsion at late stages. This growth translates into different apparent values of chirp mass in early and late stages, $m_c$ and $\tilde{m}_c$, which differ by a factor of $(1 - \tilde{\alpha}')^{2/5}$. Therefore, for models with large enough $\tilde{\alpha}'$, it will become impossible to fit the whole waveform with a single standard template with a unique chirp mass. It will be necessary to use two templates with different masses $m_E$ and $m_L$ to fit the early waveform and late waveform, respectively. The difference $m_E - m_L$ then gives the value of $\tilde{\alpha}'$, and the place where both templates fail to fit gives the mass of the dark photon. For a rigorous waveform analysis and comparison of different templates, we refer to Sampson et al. (2014); Sennett & Buonanno (2016); Sagunski et al. (2018); Shao et al. (2017); Sennett et al. (2017); Yunes et al. (2016). We leave a detailed analysis and comparison of different templates, we refer to Sampson et al. (2014); Sennett & Buonanno (2016); Sagunski et al. (2018); Shao et al. (2017); Sennett et al. (2017); Yunes et al. (2016). We leave a detailed analysis for future work.

Without going into a full waveform analysis, we outline our estimate on the size and detectability of $\tilde{\alpha}'$. From (8) it is straightforward to see that the bound is given by $\tilde{\alpha}' \lesssim \frac{1}{2}(m_{c,i} - m_{c,f})/m_c$, where $m_{c,i}$ ($m_{c,f}$) is the initial (final) chirp mass, respectively. A large enough $\tilde{\alpha}'$ corresponds to a large change of $m_c$ in the signal. We note that a rigorous bound can be drawn only if one analyzes the data using a template with $\tilde{\alpha}'$ parameter. A Markov chain Monte Carlo (MCMC) simulation Sagunski et al. (2018) using emcee Foreman-Mackey et al. (2013) shows that $|\tilde{\alpha}'| \sim 1/3$ is detectable at LIGO. This is to be compared with the “trivial” stability bound for a charged NS star, which is given by setting $E = 0$ in (7) and roughly translates to $\tilde{\alpha}' < 1$ Ray et al. (2006, 2003). Again, we note that this bound is valid only if one assumes a mediator in the range of $\mathcal{O}(10)$ km $\lesssim (m_V)^{-1} \lesssim \mathcal{O}(1000)$ km. Failure to observe a change in the chirp mass can be interpreted as either too small an $\tilde{\alpha}'$ or $m_V$ beyond the neutron star “spectrometer” window. We also note that the quantity $\tilde{\alpha}'$ is the product of the coupling strength $\alpha'$ and the charges carried by the binary $q_1 q_2$, and it is in principle possible that $\tilde{\alpha}'$ is different for different binaries due to different $q_1 q_2$. Therefore, the $\tilde{\alpha}'$ is to be measured event-wise, and thus the improvement of upper limit on $\tilde{\alpha}'$ requires a better precision of measuring $m_c$ in each event. However, we also note that for a given dark matter model, the amount of charge stored in each neutron star, and thus $\tilde{\alpha}'$, can be estimated to a certain extent, which can in turn be translated to a bound on the properties of aDM in NS. In the case of LIGO, aDM fraction can be constrained starting from the heavy end in the $m_{\chi} \sim$ TeV range. With very heavy aDM, the fraction can be constrained down to sub-percentage level. This is to be compared with 5-10% DM mass allowed in NS in models such as Foot (2004); Sandin & Ciarcelluti (2009); Fan et al. (2013), according to Ellis et al. (2017), which assumes a different DM model. We also note, the bound on $\tilde{\alpha}'$ for a specific dark matter model can be improved with increasing the statistics.

The second window comes from the GW wavelength, to which LIGO-Virgo is sensitive in the band $10^3$ km $< \lambda_{GW} < 10^4$ km. Observation of hidden sector in this window relies on a nonzero charge-mass difference $\gamma$ for the two NS. When this is the case, the dark dipole radiation is stimulated when $\lambda_{GW} \lesssim m_V^{-1}$, and at this point the waveform will develop a dipole component which is vastly different from all the other contributions. For a very light dark photon with $m_V^{-1} \gg 10^4$ km, such that $\gamma \neq 0$ in the whole observational window, the dipole component exists throughout the LIGO-Virgo band. This is illustrated in Figure 1.

We conclude by noting the following points. The existence of NS binaries itself puts a constraint on the percentage of aDM each NS carries, which is equivalent of the weakest bound in our analysis by setting $E = 0$. This sets the same bound that is required for the aDM to be contained in the NS. In addition, in this Letter we make no assumption about the structure of the dark sector or the origin of the dark boson mass. In the scenario where the dark photon mass is generated by a dark Higgs mechanism rather than a Stueckelberg mass term, the signal will likely be affected by the dark Higgs, which we leave for future work.
Similarly, we make no assumption about the mechanism of DM capture. All we assumed is that the DM mass fraction in a NS is large enough to generate an effect, while the coupled system of NS and DM remains stable. In the minimal scenario with dark repulsion the only long range force besides gravity, the stability implies that \( \dot{\alpha'} \) is always smaller than DM mass fraction. Therefore, a sensitivity of \( \dot{\alpha'} \) down to \( 10^{-2} \) means that percent-level mass fraction of DM in a NS is being constrained. We do not address how the DM got captured, but note that for instance, DM could have been captured by the adiabatically contracting gravitational potential well during the formation of the progenitor star. The above stability constraint between \( \dot{\alpha'} \) and DM mass fraction can be easily circumvented by, e.g., the presence of another force for DM which is attractive and has range \( \lesssim \mathcal{O}(10) \text{km} \).

Furthermore, if mixing between the Standard Model photon and the dark boson is assumed, the dark radiation may leave imprints in the radio frequency band. Finally, our analysis here can be readily applied to GW detectors with different frequency bands such as LISA, extending the reach of the GW spectroscopy for hidden sectors.

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