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Exclusion of the remaining mass window for primordial black holes as the dominant constituent of dark matter

Paolo Pani[1,2,*] and Abraham Loeb[2,†]

[1] *CENTRA, Departamento de Física, Instituto Superior Técnico, Universidade de Lisboa, Avenida Rovisco Pais 1, 1049 Lisboa, Portugal.*

[2] *Institute for Theory and Computation, Harvard-Smithsonian CfA, 60 Garden Street, Cambridge MA 02138, USA*

In a close encounter with a neutron star, a primordial black hole can get gravitationally captured by depositing a considerable amount of energy into nonradial stellar modes. Within a short time, the black hole is trapped inside the star and disrupts it by rapid accretion. We find that the existence of old neutron stars in regions where the dark-matter density $\rho_{\rm DM} \gtrsim 10^2(\sigma/{\rm km\ s^{-1}}){\rm GeV\ cm^{-3}}$ (where σ is the dark-matter velocity dispersion) limits the abundance of primordial black holes in the mass range $10^{17}{\rm g} \lesssim m_{\rm PBH} \lesssim 10^{24}{\rm g}$, which was previously unconstrained. In combination with existing limits, our results suggest that primordial black holes cannot be the dominant dark matter constituent.

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Primordial black holes (PBHs) might arguably be the most natural candidates to solve the dark matter (DM) puzzle. They are cold, weakly-interacting, and do not require extensions of the Standard Model of particle physics. Since being proposed by Hawking in 1971 [1] (see also [2, 3]), substantial effort has been put to detect these hypothetical objects or rule them out.

PBHs could have formed across a wide range of masses as a result of the evolving density and horizon scale of the early Universe [4, 5]. Light PBHs with a mass $m < 10^{15}$g should have already Hawking-evaporated by the present epoch [6]. On the other hand, strong constraints are in place for the DM fraction in nonevaporating PBHs with a mass $m_{\rm PBH} \gtrsim 4\times 10^{24}$g, based on a variety of dynamical [7, 8], microlensing [9–13] and other astrophysical [14–17] effects (see Ref. [4] for an overview).

Despite extensive theoretical and experimental effort, the range $10^{17}{\rm g} \lesssim m_{\rm PBH} \lesssim 10^{24}{\rm g}$ is still poorly constrained[1]. In this mass range, light PBHs could in principle explain the DM without a need for exotic, weakly-interacting particles.

Here we consider a novel phenomenon that can be used to derive stringent theoretical constraints on the range $10^{17}{\rm g} \lesssim m_{\rm PBH} \lesssim 10^{24}{\rm g}$, the only one that remains phenomenologically allowed to date. We show that in a close encounter with a neutron star (NS), light PBHs can be tidally captured by depositing a surprisingly large amount of energy into the nonradial stellar modes. Over a time scale much shorter than the NS lifetime, the PBH would be captured within the NS core and would eventually disrupt the star by rapid accretion [19, 23]. Thus, the mere observation of NSs in DM-rich environments poses a theoretical limit on the density of PBHs. Our results, in combination with previous constraints, allow to rule out PBHs as dominant DM constituent.

Tidal capture of PBHs by a NS. We consider the encounter of a PBH with mass $m_{\rm PBH}$ and a NS with mass $M \gg m_{\rm PBH}$ and radius R. For simplicity, we focus on head-on collisions, although our results can be easily generalized to arbitrary orbits. During a close encounter, the PBH will deposit energy in nonradial acoustic modes of the star. Press and Teukolsky studied an analogous process for a NS-NS encounter [24] and found that the total energy loss reads (using $G = c = 1$ units)

$$\Delta E = \frac{m_{\rm PBH}^2}{R}\sum_{l=2}^{\infty}\left(\frac{R}{R_{\rm min}}\right)^{2l+2} T_l\,, \qquad (1)$$

where $R_{\rm min}$ is the periastron distance, l is the multipolar index of the stellar modes and T_l is a function of the masses of the objects, the radius R and the periastron distance [24].

Typically, the minimum separation in a NS-NS encounter is dictated by the Roche radius within which the smallest object is tidally disrupted [25], implying $R_{\rm min} \gtrsim 3R$. Since T_l are roughly the same for any l [24], the sum in Eq. (1) converges quickly, and $\Delta E \sim m_{\rm PBH}^2/R$. However, a PBH does not get tidally disrupted; it can reach the NS surface and even travel *through* the star. One might therefore expect that the energy exchange be much larger, because higher multipoles are relevant.

To examine the role of high multipoles, we adopt an approach which is commonly used in seismology [26]. Considering a spherically symmetric, perfect-fluid star with barotropic pressure $P = P(\rho)$, and a point mass $m_{\rm PBH}$ travelling radially, we compute the seismic energy deposited into the stellar normal modes. The point-particle approximation is particularly well suited here, because in the range of interest $m_{\rm PBH}/M \lesssim 10^{-8}$. Our treatment is therefore Newtonian. When considering the motion in the interior, we neglect other effects such as accretion

* paolo.pani@tecnico.ulisboa.pt

† aloeb@cfa.harvard.edu

[1] Recent constraints include those based on femtolensing [18] and on PBHs in globular clusters [19, 20]. However, the former ignored finite-source effects and are currently being revised [21], whereas the latter are weakened by the lack of evidence for DM in globular clusters [22].

and dynamical friction [19] because, as we show, they are subdominant relative to the energy in the modes.

It is convenient to perform a normal-mode decomposition of the displacement of a fluid element [26],

$$\boldsymbol{s}(\boldsymbol{x},t) = -\mathrm{Re}\sum_{\mathfrak{n}l} \frac{c_{\mathfrak{n}l}}{\omega_{\mathfrak{n}l}^2} e^{i\omega_{\mathfrak{n}l}t} \boldsymbol{s}_{\mathfrak{n}l}\,, \tag{2}$$

where $\boldsymbol{s}_{\mathfrak{n}l}$ denotes the eigenfunction with overtone $\mathfrak{n}$ and harmonic index l and whose eigenfrequency is $\omega_{\mathfrak{n}l}$. The normal-mode equations are discussed in Ref. [26]. Here $c_{\mathfrak{n}l}$ is the excitation coefficient of the $(\mathfrak{n}, l)$ mode, obtained through a convolution of the source force density $\boldsymbol{f}(\boldsymbol{x},t)$ with the eigenfunctions:

$$c_{\mathfrak{n}l} = \int dt\, e^{-i\omega_{\mathfrak{n}l}t} \int_{\mathrm{star}} d\boldsymbol{x}^3 \partial_t \boldsymbol{f}(\boldsymbol{x},t)\cdot \boldsymbol{s}^*_{\mathfrak{n}l}(\boldsymbol{x})\,, \tag{3}$$

where the time integration is performed in the interval when the source is active. The modified seismic energy is the sum over the excitation coefficients [26]

$$\Delta E \equiv \sum_{\mathfrak{n}l} \mathcal{E}_{\mathfrak{n}l} = \frac{1}{2}\sum_{\mathfrak{n}l} \frac{|c_{\mathfrak{n}l}|^2}{\omega_{\mathfrak{n}l}^2}\,. \tag{4}$$

Once the source is specified, the knowledge of the normal modes allows us to compute $c_{\mathfrak{n}l}$ and the modified seismic energy through Eq. (4). Note that ΔE is a conservative lower bound on the total deposited energy [26].

We adopt the point-particle approximation [27]

$$\boldsymbol{f}(\boldsymbol{x},t) \equiv -\rho(\boldsymbol{x})\boldsymbol{\nabla}\Phi = m_{\mathrm{PBH}}\rho(\boldsymbol{x})\boldsymbol{\nabla}\frac{1}{|\boldsymbol{x}-\boldsymbol{x}_p(t)|}\,, \tag{5}$$

where Φ is the gravitational potential, $\boldsymbol{\nabla} = \hat{\boldsymbol{r}}\partial_r + \boldsymbol{\nabla_1}$, $\boldsymbol{\nabla_1} = \hat{\boldsymbol{\theta}}\partial_\theta + \hat{\boldsymbol{\phi}}(\sin\theta)^{-1}\partial_\phi$ and $\boldsymbol{x}_p(t) = (r_p(t), 0, 0)$ is the location of the PBH. The source term can be expanded in Legendre polynomials as

$$\boldsymbol{f}(\boldsymbol{x},t) = m_{\mathrm{PBH}}\rho(\boldsymbol{x})\boldsymbol{\nabla}\begin{cases} \sum_l \frac{r^l}{r_p^{l+1}} P_l(\cos\theta) & r_p(t) > r \\ \sum_l \frac{r_p^l}{r^{l+1}} P_l(\cos\theta) & r_p(t) < r \end{cases}$$

where we have used the fact that $\theta_p = \phi_p = 0$. This source excites only spheroidal modes, which can be conveniently decomposed in a basis of spherical harmonics,

$$\boldsymbol{s}_{\mathfrak{n}l}(\boldsymbol{x}) = U_{\mathfrak{n}l}(r)Y_{lm}(\theta,\phi)\hat{\boldsymbol{r}} + \frac{V_{\mathfrak{n}l}(r)}{\sqrt{l(l+1)}}\boldsymbol{\nabla_1}Y_{lm}(\theta,\phi)\,, \tag{6}$$

with no summation over l. Due to the symmetry of the problem, only the $m = 0$ modes are relevant and there is no explicit dependence on the coordinate ϕ.

Inserting the expansion above into Eq. (3), using Eq. (5) and the orthonormality properties of the spherical harmonics, we can compute $c_{\mathfrak{n}l}$. For the external motion we obtain

$$c^{\mathrm{ext}}_{\mathfrak{n}l} = -m_{\mathrm{PBH}}\int_{-\infty}^{t_R} dt \frac{l(l+1)v_p(t)}{r_p^{l+2}(t)} e^{-i\omega_{\mathfrak{n}l}t} \times \int_0^R dr\rho(r)r^{l+1}\left(U_{\mathfrak{n}l} + \frac{\sqrt{l(l+1)}}{l}V_{\mathfrak{n}l}\right)\,, \tag{7}$$

where $r_p(t < t_R) = \left(R^{3/2} + 3\sqrt{M/2}(t_R - t)\right)^{2/3}$, $v_p = \partial_t r_p$, and t_R marks the time at which the particle reaches the surface of the star. Conservatively, we assume that the PBH is at rest at infinity.

The response when the PBH travels within the NS is more involved and it has to be divided in the case $r_p(t) < r$ and $r_p(t) > r$. The final result reads

$$\begin{aligned} c^{\mathrm{int}}_{\mathfrak{n}l} &= -m_{\mathrm{PBH}}\int_{t_R}^{t_0} dt \frac{l(l+1)v_p(t)}{r_p^{l+2}(t)} e^{-i\omega_{\mathfrak{n}l}t} \\ &\times \int_0^{r_p(t)} dr\rho(r)r^{l+1}\left(U_{\mathfrak{n}l} + \frac{\sqrt{l(l+1)}}{l}V_{\mathfrak{n}l}\right) \\ &-m_{\mathrm{PBH}}\int_{t_R}^{t_0} dt\, l(l+1)v_p(t)r_p^{l-1}(t)e^{-i\omega_{\mathfrak{n}l}t} \\ &\times \int_{r_p(t)}^{R} dr\frac{\rho(r)}{r^l}\left(U_{\mathfrak{n}l} - \frac{\sqrt{l(l+1)}}{l+1}V_{\mathfrak{n}l}\right)\,, \end{aligned} \tag{8}$$

where $t_0 > t_R$ is the time the particle reaches the center of the star and $r_p(t > t_R)$ is obtained solving the equation of motion of a point particle travelling through a density distribution $\rho(r)$. Finally, the seismic energy is computed through Eq. (4) with $c_{\mathfrak{n}l} = c^{\mathrm{int}}_{\mathfrak{n}l} + c^{\mathrm{ext}}_{\mathfrak{n}l}$, and multiplying the result by a factor 2 to account for the symmetric motion from the center to the exterior of the NS.

We compute $\mathcal{E}_{\mathfrak{n}l}$ numerically up to $l = 100$ for a variety of NS equations of state (EOS). The contribution of the overtones ($\mathfrak{n} > 0$) is negligible and in practice we consider only the fundamental $\mathfrak{n} = 0$ modes. The multipolar contributions as functions of l are shown in Fig. 1 for typical configurations.

Our analysis reveals a remarkable property which, to the best of our knowledge, has never been discussed before. For very compact objects, the multipolar contri-

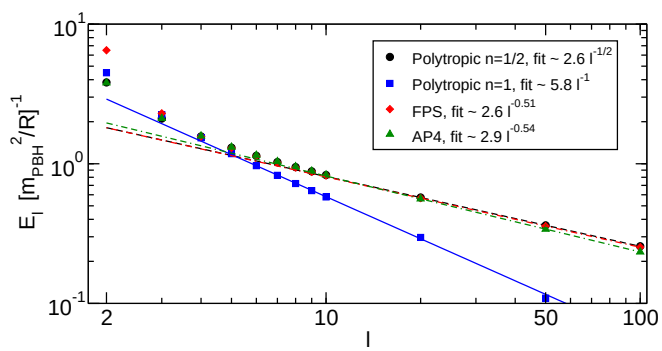

FIG. 1. Multipolar contributions $\mathcal{E}_{\mathfrak{n}l}$ to the total energy $\Delta E = \sum_{\mathfrak{n}l}\mathcal{E}_{\mathfrak{n}l}$ as a function of l and for the dominant $\mathfrak{n} = 0$ terms. For polytropic stars with $P(\rho) = k\rho^{1+1/n}$ our results imply $\mathcal{E}_{\mathfrak{n}l} \propto l^{-n}$ in the large-l limit. For the tabulated FPS and AP4 EOS the energy scales as $\mathcal{E}_{\mathfrak{n}l} \propto l^{-0.51}$ and $\mathcal{E}_{\mathfrak{n}l} \propto l^{-0.54}$, respectively. For all models the NS radius $R \approx 12$km, and the mass $M \approx 2M_\odot$ and $M \approx 1.4M_\odot$ for the polytropic models with $n = 1/2$ and $n = 1$, respectively, and $M \approx 1.5M_\odot$ and $M \approx 1.9M_\odot$ for the FPS and AP4 EOS, respectively.

bution $\mathcal{E}_{\mathfrak{n}l}$ decays very slowly in the large-l limit. For a polytropic EOS, $P(\rho) = k\rho^{1+1/n}$, we find the following behavior in the large-l limit,

$$\mathcal{E}_{\mathfrak{n}l} \sim \frac{m_{\rm PBH}^2}{R}\frac{\gamma}{l^n} \qquad l \gg 1\,, \tag{9}$$

where the constant γ is of order unity and its precise value depends on the EOS. Although Fig. 1 shows only the cases $n = 1$ and $n = 1/2$, we verified Eq. (9) for $0 \leq n \leq 3$. This result is not an artifact of the polytropic-fluid approximation. Indeed, we have considered two representative tabulated EOS, namely FPS [28] and AP4 [29], obtaining a result analogous to Eq. (9) but with $n \approx 0.51$ and $n \approx 0.54$, respectively. This is natural because NS equilibrium structures are roughly described by polytropes with $0.5 < n < 1$.

Equation (9) implies that the total energy loss, $\Delta E = \sum_{\mathfrak{n}l} \mathcal{E}_{\mathfrak{n}l}$, *diverges* for polytropes with $n \leq 1$ and for realistic NS models. A similar phenomenon occurs in the head-on collision of a point-particle with a BH. In that case, the energy *absorbed* at the horizon diverges [30]. Our results show that a similar divergence also occurs if the central object is a NS. This divergence is enabled by the fact that a PBH does not suffer tidal disruption; instead it reaches the radius of the star and even travels within it. In fact, an analysis of the excitation factors (7) and (8) reveals that the majority of the energy is released when the particle reaches the surface and when it passes through the outer layers of the NS, where the large-l eigenfunctions peak. To compute the eigenfunctions in the NS crust region precisely, tabulated EOS were also implemented using piecewise functions [31]. Although during the inside motion the energy loss is roughly an order of magnitude larger than during the outside motion, the divergence of ΔE would also occur if only the outside motion is taken into account, provided the PBH can reach the surface without being destroyed.

For an incompressible fluid with $n = 0$, Eq. (9) predicts $\mathcal{E}_{\mathfrak{n}l} \approx$ const in the large-l limit, precisely as in the BH case [30]. The $n = 0$ case can be solved analytically, confirming that the divergence of the total energy is not a numerical artifact.

A finite-size cutoff for the point particle can regularize such divergence. In practice, the multipolar sum (4) should be truncated at $l_{\rm max} \sim \frac{\pi}{2}\left(\frac{R}{2m_{\rm PBH}}\right) \gg 1$ to account for finite-size effects on the scale of the PBH horizon [30]. The total energy is then finite and reads

$$\Delta E = \frac{m_{\rm PBH}^2}{R}\frac{2\gamma}{(1-n)}\left(\frac{\pi}{4}\right)^{1-n}\left(\frac{R}{m_{\rm PBH}}\right)^{1-n}, \tag{10}$$

where we assumed $n < 1$. Note that this energy release is larger than the energy associated with accretion and dynamical friction [19] by a factor $(R/m_{\rm PBH})^{1-n} \gg 1$. Nonetheless, it is much smaller than the PBH rest mass, consistently with our perturbative approach. We stress that Eq. (10) applies also to realistic EOS which effectively correspond to $n \approx 0.5$.

Tidal capture rate. During a collision with a NS, a PBH loses an amount of energy given by Eq. (10), and might get gravitationally captured by the star. The time scale for this process and the associated capture rate were computed in Ref. [19]. Given a radial orbit with apastron $r_{\rm max}$, the half period of the orbit reads $\Delta T = \pi r_{\rm max}^{3/2}/\sqrt{M}$. Due to the energy loss, $dE_p/dt \sim -\Delta E/\Delta T$, the apastron decreases at each encounter. Eventually, the PBH would be confined in the NS core in a time scale defined by $r_{\rm max}(t_{\rm loss}) = 0$, i.e. $t_{\rm loss} \sim 2\pi m_{\rm PBH}\sqrt{M r_{\rm max}^0}/\Delta E$, where $r_{\rm max}^0 = r_{\rm max}(t = 0)$ can be estimated by requiring that the initial energy is of the order of ΔE, i.e. $r_{\rm max}^0 \sim M m_{\rm PBH}/\Delta E$. Based on Eq. (10), this time scale depends on the NS configuration. For $n = 1/2$, we obtain that only PBHs with

$$m_{\rm PBH} \gtrsim 3\times 10^5 \left(\frac{M}{1.4M_\odot}\right)^{4/3}\left(\frac{R}{12{\rm km}}\right){\rm g}\,, \tag{11}$$

can be captured on a time shorter than a typical NS lifetime, $t_{\rm NS} \sim 10^{10}$yr. Qualitatively similar results hold for other values of n in Eq. (10). The bound (11) shows that any nonevaporating PBH is captured on a time shorter than the NS lifetime.

The estimate above ignores possible interactions between the modes excited in a past passage and the motion of the PBH during the next passage. Resonances might occur when the orbital frequency ΔT^{-1} is comparable to the mode frequencies which, for the fundamental f-modes, scale as $\omega_{0l}^2 \sim 2l(l+1)/(2l+1)M/R^3$ (see e.g. Ref. [32]). This frequency is comparable to ΔT^{-1} when $l = \mathcal{O}(1)$ and $r_{\rm max} \sim R$, i.e. for low multipoles and only when the PBH is already confined within the NS. Our analysis is conservative, because it ignores possible resonances for which the energy loss would be amplified when $r_{\rm max} \lesssim R$, reducing the time scale for tidal capture.

Let us now estimate the capture rate. Assuming that each PBH – with specific energy E_p and specific angular momentum J_p – follows a Maxwellian velocity distribution with dispersion σ, the capture rate can be obtained by considering the region of the (E_p, J_p) parameter space corresponding to those orbits whose periastron is at most equal to the radius of the star and that lose enough energy to become gravitationally bound after the first encounter. The latter condition is simply $E_p < \Delta E$. The former condition can be obtained by analyzing the geodesic motion of a nonrelativistic particle. The final result for the capture rate is [33]

$$F = \sqrt{6\pi}\frac{\rho_{\rm PBH}}{m_{\rm PBH}}\frac{2MR}{\sigma(1-2M/R)}\left[1 - e^{-\frac{3\Delta E}{m_{\rm PBH}\sigma^2}}\right]. \tag{12}$$

where the mass density of PBHs can be defined in terms of the DM density through $\rho_{\rm PBH} = f_{\rm PBH}\rho_{\rm DM}$, with $f_{\rm PBH}$ being the DM fraction in PBHs.

After being captured, a PBH with a mass satisfying Eq. (11) would reach the center of the star in a time $t_{\rm loss} \ll t_{\rm NS}$ and would quickly disrupt the star by rapid accretion [23]. Therefore, the mere existence of old NSs

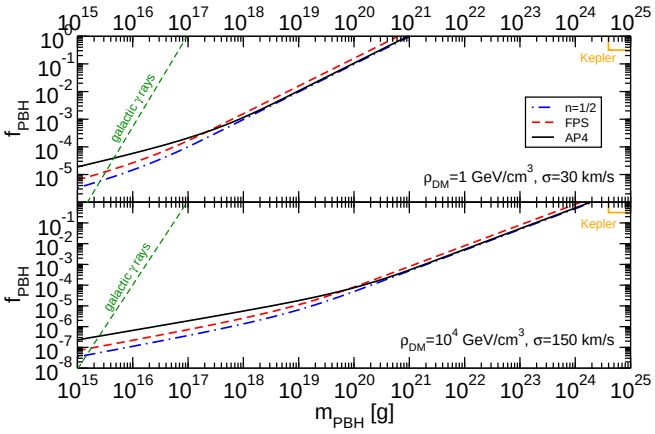

FIG. 2. Constraints on the DM fraction in PBHs $f_{\rm PBH}$, derived from PBH capture by a NS in the PBH mass range 10^{15}g $\lesssim m_{\rm PBH} \lesssim 10^{25}$g currently unconstrained by observations. We consider a polytropic EOS with $n = 1/2$ and the FPS and AP4 tabulated EOS with the same parameters as in Fig. 1. We consider typical values of the DM density $\rho_{\rm DM}$ and velocity dispersion σ for the centers of the Large Magellanic Cloud [34] (top panel) and the Milky Way [35] (bottom panel). The orange rightmost top curve shows the existing limit from Kepler [13], whereas the green leftmost line is the limit imposed by observations of the extragalactic photon background [4].

in DM-rich galaxies implies that the survival probability, $\exp(-t_{\rm NS}F)$, is not small. This translates to an upper bound on the DM fraction in PBHs:

$$f_{\rm PBH} < \frac{m_{\rm PBH}\sigma(1-2M/R)}{2\sqrt{6\pi}t_{\rm NS}MR\rho_{\rm DM}}\left[1-e^{-\frac{3\Delta E}{m_{\rm PBH}\sigma^2}}\right]^{-1} . \quad (13)$$

Limits on DM fraction in PBHs. Given a detection of an old NS in a DM-rich environment, Eq. (13) can be used to put a constraint on the DM fraction in PBHs. This constraint is inversely proportional to the DM density $\rho_{\rm DM}$. Because slow PBHs are captured more easily, the most stringent constraints can be derived from galaxies with a low velocity dispersion and a high DM density. In Fig. 2 we show the exclusion plots for two representative examples: a region of moderate local DM density, $\rho_{\rm DM} \sim 1$GeV cm^{-3}, with low velocity dispersion $\sigma \sim 30$km s^{-1}, and a region of high DM density, $\rho_{\rm DM} \sim 10^4$GeV cm^{-3}, but with high velocities, $\sigma \sim 150$km s^{-1}. The former choice is compatible with observations at about 1kpc from the galactic center of the the Large Magellanic Cloud [34], whereas the latter corresponds to the center of the Milky Way in the conservative case of a mild DM spike [35]. NSs are known to exist in the Large Magellanic Cloud [36, 37] as well as near the Milky Way center [38, 39].

Discussion. Under normal circumstances, the energy deposited in the modes of a star during a close encounter with a PBH would be $\Delta E \sim m_{\rm PBH}^2/R$ [24], at most comparable to other effects such as energy loss due to accretion and dynamical friction. We have found that when a PBH travels within a NS, the energy loss is much higher and given by Eq. (10). This has profound implications: nonevaporating PBH can be trapped within the NS core in a time scale shorter than the NS lifetime and, if such capture occurs, the NS would quickly be destroyed through rapid accretion of its nuclear material onto the PBH. Thus, observations of NSs in DM-rich environments can be used to derive stringent constraints on the DM fraction in PBHs.

The absence of detection of microlensing events in the Kepler data – together with previous bounds [4] – sets the lower limit $m_{\rm PBH} \gtrsim 4\times10^{24}$g [13]. An analysis of Eq. (13) shows that any NS detection in a region where the local DM density satisfies

$$\rho_{\rm DM} \gtrsim 240\left(\frac{\sigma}{\rm km\ s^{-1}}\right){\rm GeV\ cm^{-3}}\,, \quad (14)$$

places new theoretical constraints on the DM fraction in nonevaporating PBHs with mass $m_{\rm PBH} \lesssim 4\times10^{24}$g. In deriving Eq. (14) we assumed a NS with $M \sim 1.4M_\odot$ and $R \sim 12$km, but the prefactor would be of the same order for different choices of M and R.

As shown in Fig. 2, PBHs with $m_{\rm PBH} \lesssim 10^{17}$g are already excluded by observations of the extragalactic photon background [4]. PBHs with $m_{\rm PBH} \gtrsim 10^{17}$g are larger than the size of a neutron by a factor of $\sim 10^2$, implying that the fluid approximation is valid for matter with density as low as $\sim 10^9$g cm^{-3}, i.e. even in the NS crust. Therefore, our analysis excludes a range of more than seven orders of magnitude, 10^{17}g $\lesssim m_{\rm PBH} \lesssim 10^{24}$g, which was previously viable. In combination with previous bounds, this suggests that PBHs of any mass cannot be the dominant constituent of the DM.

Our theoretical bounds are only mildly dependent on the NS EOS and they are more sensitive to macroscopic properties such as the NS mass and radius. This is because they crucially depend on the energy deposited in the high-l modes, which are mostly localized near the NS crust, where the behavior of matter is well known and all realistic EOS are equivalent.

Our Newtonian approach provides a *conservative* estimate, because it does not account for the energy loss due to gravitational-wave emission during the encounter, nor for the energy deposited in relativistic w-modes. Old NSs rotate rather slowly and the order of magnitude of our bounds should remain valid also in the spinning case [40].

Finally, we note that there are currently no NS or pulsar detections in galaxies with very high DM densities and low velocity dispersion, such as the Milky Way's dwarf spheroidal satellites [41]. NS discoveries in these systems would place much tighter constraints than those presented here.

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