Positrons from Primordial Black Hole Microquasars and Gamma-ray Bursts

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We present several scenarios how capture of primordial black holes (PBHs) with masses $10^{-16} M_\odot \lesssim M_{\text{PBH}} \lesssim 10^{-7} M_\odot$ by compact objects (white dwarfs or neutron stars) can source short gamma-ray bursts (sGRBs) as well as microquasars (MQs). A small PBH captured by a compact star will eventually consume the host, turning it into a stellar-mass BH. We argue that formation of an accretion disk around the resulting BH, which is an important prerequisite for standard sGRB production mechanisms, can be generic. The PBH-induced relativistic sGRB flares as well as continuous MQ jets will accelerate positrons to high energies. We find that if PBHs constitute a few percent or more of the dark matter, the generated positrons can address the excess observed in the positron flux by the Pamela, the AMS-02 and the Fermi-LAT experiments. The allowed resulting parameter space for primordial black holes to constitute dark matter is in a favorable region to permit for resolution of several other astronomical puzzles.

Introduction — Primordial black holes (PBHs) can appear from early Universe dynamics and account for all or part of the dark matter (DM) [1–13]. Aside from being a theoretical curiosity, null search results for conventional DM particle candidates [14–15] as well as the possible implications [16–23] for the newly-opened field of gravity-wave astronomy [24–26] further elevate the interest in PBH-related studies. Recent investigations have shown that PBHs could shed light on a variety of outstanding astronomical puzzles, such as the origin of $r$-process nucleosynthesis material, fast radio bursts and the 511-keV line [27].

Observations by several experiments, including PAMELA [28], Fermi-LAT [29] and AMS-02 [30], have identified a rise in positron cosmic ray flux above $\sim 30$ GeV. The origin of this phenomenon remains elusive. A multitude of proposals have been put forward to address it, which can be generally grouped together as those based on astrophysical sources (e.g. pulsars [31–36], supernova remnants [37–41], microquasars [42]), particle dark matter annihilations/decays [43–46] as well as those based on cosmic ray propagation effects [47–49].

In this Letter we discuss how PBHs can incite gamma-ray bursts (GRBs) and participate in the formation of microquasars (MQs), either of which has the capacity for accelerating particles to high energies and, as we demonstrate, provide an explanation of the positron excess. Heuristically, if a small PBH with mass $10^{-16} M_\odot \lesssim M_{\text{PBH}} \lesssim 10^{-7} M_\odot$ is captured by a compact object, a white dwarf (WD) or a neutron star (NS), it will eventually consume the host (e.g. [50]) and result in a stellar-mass BH. The system’s energy, released on dynamical time-scales, is sufficient to power a short GRB. In a different scheme, the resulting stellar-mass BH could also steadily accrete matter if the considered star system is a binary, thus powering a microquasar jet. While a GRB explosion provides a singular injection of a large quantity of energy, a microquasar jet is a continuous energy injection source. Relativistic positrons, accelerated either through a burst or a jet, will diffuse and are observable. Our proposal to address the detected positron excess thus unifies the beneficial features of explanations based on astrophysical sources with dark matter.

Black hole capture — A small PBH can become gravitationally captured by a NS or a WD if it loses sufficient energy through dynamical friction and accretion as it passes through the star (see Figure 1, left). We briefly review the main capture ingredients, following [23, 50]. The full capture rate is given by $F = \langle (\Omega_{\text{PBH}}/\Omega_{\text{DM}}) F_0 \rangle$, where $\Omega_{\text{PBH}}$ is the PBH contribution to the overall DM abundance $\Omega_{\text{DM}}$. The base Galactic capture rate $F_0$ is

$$F_0 = \sqrt{6\pi} \frac{\rho_{\text{DM}}}{M_{\text{PBH}}} \left[ \frac{R_{\text{NS}}R_{\text{b}}}{\tau(1 - R_{\text{b}}/R_{\text{NS}})} \right] \left(1 - e^{-E_{\text{loss}}/E_b}\right),$$

where $\rho_{\text{DM}}$ is the DM density, $M_{\text{PBH}}$ is the PBH mass, $\tau$ is the DM velocity dispersion (assumed to have Maxwellian distribution), $E_b = M_{\text{PBH}}v^2/3$, $R_{\text{NS}}$ is the radius of the NS with mass $M_{\text{NS}}$ and Schwarzschild radius $R_S = 2G M_{\text{NS}}$. If the interaction energy loss $E_{\text{loss}}$ exceeds kinetic energy of the PBH, then it will be captured. The average energy loss for a NS is $E_{\text{loss}} \approx 58.8 G^2 M_{\text{PBH}}^2 M_{\text{NS}}/R_{\text{NS}}^2$. Throughout this work we consider a typical NS to have radius $R_{\text{NS}} \sim 12$ km, mass $M_{\text{NS}} \sim 1.5 M_\odot$ and spinning with a milli-second period $P \sim 1$ ms (i.e. milli-second pulsar) at an angular velocity $\Omega_{\text{NS}} = 2\pi/P$. In the case of WDs the star’s mass is some-what lower, but the radius is significantly larger. This results in a WD capture rate being several orders below that of a NS. The total number of PBHs captured within time $t$ is given by $F t$.

After capture, PBH will settle inside the star and consume it through Bondi spherical accretion. For a typical NS, the time for captured PBH to settle is $t_{\text{set}}^{\text{NS}} \approx 9.5 \times 10^4 (M_{\text{PBH}}/10^{-11} M_\odot)^{-3/2}$ yrs. For a WD it is $t_{\text{set}}^{\text{WD}} \approx 6.4 \times 10^6 (M_{\text{PBH}}/10^{-11} M_\odot)^{-3/2}$ yrs. Once set-
tled inside, the time for black hole to consume the star is \( t_{\text{con}}^{\text{NS}} \approx 5.3 \times 10^{-3} (10^{-11} M_\odot / M_{\text{PBH}}) \) yrs for a NS and \( t_{\text{con}}^{\text{WD}} \approx 2.9 \times 10^{-2} (10^{-11} M_\odot / M_{\text{PBH}}) \) yrs for a WD, respectively. If interaction time exceeds the timescales associated with the above processes, the system will effectively contain a stellar-mass BH once a PBH has been captured. We note that this outcome does not strongly depend on the star’s equation of state (EoS).

**Gamma-ray bursts** — Short GRBs \([51, 52]\) are irregular electromagnetic pulses that last \( \sim 0.1 - 2 \) s with a total \( \gamma \)-ray energy release of \( \sim 10^{48} - 10^{50} \) erg. The prevalent progenitor scenario for their origin is binary compact object mergers, NS-NS \([53]\) or BH-NS \([54]\). This picture is further supported by recent observation of a NS-NS merger with accompanying electromagnetic emission \([55]\). Another possibility is an accretion-induced collapse of NSs \([56, 58]\). Long GRBs (timescale > 2 s) are thought to originate from “failed supernovae” of massive stars (i.e. “collapsars”) \([59]\). A common theme in the above scenarios is a resulting BH that is expected to be engulfed in debris forming a surrounding accretion disk/torus. The disk is rapidly accreted on dynamical timescales and the system releases energy as a GRB. As we argue below, disk formation and the resulting BH-disk system is also a likely consequence of PBH-induced NS implosions. Further, we outline several generic channels of sGRB production for this setup. We note that while PBH-induced GRBs have already been mentioned in \([60]\), their scenario is based on model-dependent NS quark-plasma phase transitions rather than an accretion disk and requires existence of hypothetical stable quark star configurations.

A sizable accretion disk of mass \( M_d / M_\odot \sim 10^{-3} - 10^{-1} \), where \( M_\odot \) denotes the total mass-energy, can generically form during collapse of a uniformly rotating star that is spinning near mass-shedding limit \([61]\). The size of the disk, however, strongly depends on the initial conditions of the system and the star’s equation of state. The EoS can be modeled by a polytropic relation \( P = K \rho^{1+1/n} \), where \( P \) is the pressure, \( K \) is a constant, \( \rho \) is the density and \( n \) is the polytropic index. For NSs, recent observations of \( \sim 2 M_\odot \) pulsars J1614-2230 \([62]\) and J0348+0432 \([63]\) favor a stiffer EoS, corresponding to a smaller polytropic index and a more uniform star density profile. In \([64]\), however, it was shown that for stiffer EoS formation of an accretion disk from a collapsing rotating massive NS is problematic. This pitfall can be circumvented in our setup with a PBH. Heuristically, this can be seen by noting that with an extremely dense PBH growing at the center of the star the system effectively resembles a softer (more cored) EoS, even if the NS envelope surrounding the BH has a stiff profile.

In a recent study \([27]\) we have shown that up to \( \sim 0.1 - 0.5 M_\odot \) of neutron-rich material can be ejected from a NS that is rotating near mass-shedding limit if all of the angular momentum from particles falling into the BH is efficiently transferred to the star’s outer shells and the star remains a rigid rotator throughout the BH evolution. On the other hand, it is likely that some of the angular momentum will be transferred to the BH instead. Here, we consider that the BH acquires majority of the original star’s angular momentum. This will result in an increase of the BH’s Kerr spin parameter \( a \), which in turn decreases its innermost stable orbit (ISCO) radius \( r_{\text{ISCO}} \). For a BH of mass \( M_{\text{BH}} \), the ISCO radius can vary from \( r_{\text{ISCO}}(a = 0) = 3 R_\odot \) for a non-spinning Schwarzschild BH to \( r_{\text{ISCO}}(a = 1) = R_\odot / 2 \) for a maximally spinning Kerr BH, with \( R_\odot = 2 G M_{\text{BH}} \) denoting the BH’s Schwarzschild radius. Any residual material outside of \( r_{\text{ISCO}} \) will participate in formation of the accretion disk. We estimate that a \( \sim 10^{-2} - 10^{-1} M_\odot \) disk formation occurs in a generic NS-PBH system (see Supplemental Material \([65]\)). Upcoming simulations will allow to definitively verify this claim. As a base reference for estimates below we take a typical disk surrounding the resulting BH to be of \( M_d \approx 0.1 M_\odot \) size and further consider it to be accreted within \( \Delta t \sim 0.1 \) s. The duration of the disk’s accretion sets the burst emission timescale. As alluded to above, the sGRB jet engine \([52]\) can be powered by extracting energy from the accretion disk. For a maximally rotating BH up to 42% of the disk’s binding energy \( E_b = 0.42 M_d \approx 10^{53} (M_d/0.1M_\odot) \) erg is extractable. The two major production mechanisms are neutrino–anti-neutrino annihilation and magnetohydrodynamic (MHD) winds (Blandford-Payne \([66]\)). Unlike the post-merger NS and the collapsar scenarios, in our setup there is no significant source of potential energy that will allow for the resulting accretion disk to be sufficiently heated. Hence, we do not envision that neutrino–anti-neutrino production will play an important role. On the other hand, shearing induced by differential rotation of the disk as well as instabilities \([67]\) are expected to amplify the disk’s magnetic fields to magnetar-like \( \sim 10^{15} \) G levels (e.g. \([68]\)). If the magnetic field does not penetrate the BH and is confined to the disk, the disk will radiate as an electro-magnetic dipole similar to a pulsar (see model of \([69]\)) with a resulting luminosity of \( L_{\text{EM}} \sim 10^{50} (B/10^{15} \text{G})^2 (P/\text{1ms})^{-4} (R/10 \text{km})^6 \) erg s\(^{-1}\) for a solar mass Schwarzschild BH with \( R = r_{\text{ISCO}}(a = 0) \).

An alternative sGRB power source is the BH itself. After PBH-NS interaction, the remaining solar mass BH can harbor significant rotational energy. If the disk’s magnetic field penetrates the BH, then rotational energy can be extracted from the BH’s ergosphere through the Blandford-Znajek (BZ) mechanism \([70]\) as depicted on Figure \([1]\) (right). For a maximally spinning BH up to 29% of its rest mass is extractable \([71]\). The total NS angular momentum is given by \( J_{\text{NS}} = \Omega_{\text{NS}} I_{\text{NS}} \), where \( I_{\text{NS}} = (2/5) M_{\text{NS}} R_{\text{NS}}^2 \) is its moment of inertia. For a BH retaining most of the original NS mass as well as angular momentum, the resulting Kerr parameter will be \( a = J_{\text{BH}} / G M_{\text{BH}}^2 \approx J_{\text{NS}} / G M_{\text{NS}}^2 \approx 0.53 (P/\text{1ms})^{-1} \). The
available energy for extraction is the reducible BH mass, given by

\[ M_{\text{red}} = M_{\text{BH}} \left[ 1 - \left( \frac{f(a)}{2} \right)^{1/2} \right] \approx 10^{53} \text{ erg} , \quad (2) \]

where \( f(a) = (1 + \sqrt{1 - a^2}) \). The frame-dragging BH angular velocity at the outer event horizon \( R_+ = f(a)R_S/2 \) is \( \Omega_H = a/2R_+ \approx 2 \times 10^4 \text{ rad s}^{-1} \). For a uniform magnetic field \( B_0 \) aligned with rotation axis of the BH in vacuum (i.e. Wald solution [72]) the flux through BH’s hemisphere is given by [71]

\[ \Phi_{\text{BH}} = \pi R_+^2 B_0 \left[ 1 - a^4 \left( \frac{R_S}{2R_+} \right)^4 \right] . \quad (3) \]

The resulting BZ luminosity is thus [70, 73]

\[ L_{\text{BZ}} \approx \frac{\kappa}{4\pi} \Omega_H^2 \Phi_{\text{BH}}^2 = 6 \times 10^{47} \left( \frac{B_0}{10^{15} \text{ G}} \right)^2 \text{ erg s}^{-1} , \quad (4) \]

where constant \( \kappa \approx 0.05 \) weakly depends on the magnetic field geometry. The above approximation is accurate up to \( a \approx 0.95 \) [73], which covers our region of interest. Hence, the total energy output of the process is \( E_{\text{BZ}} = L_{\text{BZ}} \Delta t \approx 10^{47} \text{ erg} \).

A more speculative recent proposal of [74] could lead to yet another GRB production channel. The authors argue (however, see [75]) that if magnetosphere is properly accounted for, the BH can in fact retain “magnetic hair” for a long time even without any supporting disk. In that case, the BH itself can source the BZ jet on timescales exceeding \( \Delta t \sim 1 \text{ s} \), resulting in a very efficient GRB production site. These claims strongly depend on details of magnetic reconnection that are highly uncertain.

We note that magnetic fields of \( \sim 10^{16} \text{ G} \) strength can already be found in magnetars [73, 77], which themselves have been proposed as sources of GRBs. While the magnetar population could be significant, magnetic braking and very rapid radiative spin-down from millisecond star period at birth to a period of \( O(1 - 10) \text{ s} \) renders their contribution to PBH-star interactions that we are interested in as insignificant.

**Microquasars** — Having discussed GRB production we now turn to microquasars [78, 79]. They are X-ray binaries (XRBs) with accreting stellar-mass compact objects (NS or BH) that exhibit relativistic jets. Broad MQ emission spectrum [80] spans many decades in energy, ranging from eV up to and exceeding TeV. PBHs can also form MQs, with the most straightforward possibility being companion capture by a stellar-mass PBH. While such PBH mass range is interesting in light of recently detected GW signals [16], it is already strongly constrained. Here we envision a different scenario, which is realized when a small PBH consumes a binary star (NS or WD) and transforms it into a jet-emitting BH.

We first consider the case of forming a BH MQ from a binary NS that interacted with a PBH. Even before turning into BHs, neutron star XRBs are likely already emitting relativistic jets. Still, the resulting behavioral change between BH and NS MQs is observable. One discriminating signature between the two is the emission of keV thermal X-ray component originating from the plasma-star surface interactions, which is present in the spectrum of XRBs with NSs but not with BHs since the latter have no star surface [81]. In principle, the resulting BH MQs might also be significantly more efficient emitters. The associated BZ mechanism that allows to harvest energy from a rotating BH is known to be able to extract energy at a rate higher than the energy inflow to the BH, thus exceeding 100% efficiency, as occurs in the...
case of magnetically-arrested disks [52]. Observations of stellar-mass black hole binaries, however, appear to not indicate a significant spin-jet correlation for continuous (hard MQ state) jets [53], thus weakening support for a BZ mechanism at play. We note that a potential spin-jet correlation has been reported for transient jets [54]. For further comparison of BH vs. NS XRBs see [55].

Interacting white dwarfs also possess accretion disks and can emit jets. While there have been proposals to model them after MQs, as “nano-quasars” [56], their jets are non-relativistic ($v \sim 10^{-2}$) and luminosity is strongly sub-Eddington ($L \lesssim 10^{-2} L_{\text{Edd}}$). This picture can dramatically change if a PBH transforms an accreting WD into a BH. The resulting BH mass will be similar to the original WD, $M_{\text{WD}} \sim 1 M_\odot$. The accretion distance, however, will drastically decreases from the WD radius of $R_{\text{WD}} \sim 10^3 - 10^4$ km to the BH ISCO radius of $\gamma_{\text{ISCO}} < 10$ km. For a constant accretion rate $\dot{M}$ the respective luminosity is given by $L_{\dot{M}} = \frac{\dot{M} M c^2}{1 + \epsilon}$, where $\epsilon = GM/R$ is the efficiency parameter [81]. Hence, the accretion efficiency of the system increases from $\sim 10^{-4}$ for the WD to $\sim 10^{-1}$ for the BH, which could enable luminosity to reach near-Eddington levels as desired for microquasar jets [80]. Since a typical WD rotates with a period of hours-to-days [87], the spin parameter of the (post-WD) resulting BH is negligible. Hence, activation of a jet through BZ is not expected to be effective and other mechanisms will need to be invoked. For a Blandford-Payne-type setup the jet luminosity will depend on the accretion as $L_{\text{jet}} = (1/2) q_{\text{jet}} L_{\text{acc}}$, where $q_{\text{jet}} < 1$. We note that the decrease in the compact object radius will also result in an increase of the temperature in the surrounding Shakura-Sunyaev thin accretion disk [88]. The maximum temperature scales with the disk’s inner radius $r_0$ as $T_{\text{max}} \propto (M M/\dot{M}^2)^{1/4}$. For a constant accretion rate, active galactic nuclei-MQ jets are known to obey general scaling relations with size of the central black hole [79, 89]. Hence, our MQ with a solar mass BH could in principle produce similar output as a typical MQ with BH of a some-what higher mass.

**Positrons** — Both GRBs as well MQ jets can accelerate particles to high energies. First we comment on GRBs. One way to ensure that a jet will possess a high relativistic factor $\Gamma$ is by powering the GRB through the BZ mechanism. High magnetic fields will prevent jet contamination by protons, which will enable jet to be baryon-dilute and hence efficiently accelerate lighter particles such as electrons and positrons. GRBs are expected to be sources of copious electron-positron pair production at the MeV energies [90, 91]. A sizable population of GeV-TeV positrons can appear from re-scattering of TeV photons from the initial burst on a low energy target photon field, resulting in $\gamma \gamma \rightarrow e^{+}e^{-}$. Such field can be provided by $eV$ photons from the optical flash afterglow associated with the GRB in question [92]. Assuming typical GRB energetics, the resulting total pair energy will be of similar magnitude as the GRB photons, carrying $\sim 10^{50}$ erg and with a spectral index $\sim 2$ [93].

For MQs, a variety of jet models exist [80]. Their jets include several contributions, including the corona, the accretion disk and the companion star fields, resulting in a complicated spectrum with photon, lepton and hadron components. A major production channel (e.g. see leptohadronic models of [94]) for energetic GeV-TeV positrons is $p \gamma \rightarrow \pi^+ n, \pi^0 p$ followed by pion decays to leptons. While other production channels such as $p \gamma \rightarrow p e^+ e^-$ exist, they are typically suppressed or yield positrons with lower energy. The resulting emission is well modeled by a broken power-law with a spectral index of $\sim 1.5 - 2.5$ [95]. Energetic neutrinos produced from MQ muon and pion decays have also been suggested as signal sources for neutrino detectors such as IceCube [96].

The emitted positrons will diffuse, with evolution characterized by the diffusion equation. For a burst point source with a power-law spectrum, such as a GRB, the diffusion equation can be fully solved analytically (see e.g. [95]). In [95] the authors have fitted the observed positron excess to the diffused MQ spectrum, assuming a uniform Galactic distribution of XRBs that are currently emitting. They found that $O(10^{-2} - 10^{-3})$ MQs with a typical Eddington luminosity of $\sim 10^{38}$ erg s$^{-1}$ and emitting $\sim 10^{32} - 34$ erg s$^{-1}$ positron flux can account for the ob-
served excess. Diffusing positrons will lose energy primarily through synchrotron and inverse Compton scattering. The energy loss is described by $\frac{dE}{dt} = -\beta E^2$, where $\beta \simeq 10^{-16}$ GeV$^{-1}$ s$^{-1}$ \cite{103}. Hence, the diffusion time to lose half of initial energy $E_0$ is $\tau = 1/\beta E_0$, which for a 100 GeV positron translates to $\tau (100 \text{ GeV}) \simeq 3 \times 10^6$ yrs. In \cite{23} it was shown that a single GRB burst that occurred at a distance $\sim 1$ kpc away within the diffusion time, or conversely a single MQ that has been continuously emitting for such duration, can both fit well the observed excess.

### Signal rate

We now comment on the contribution to the positron excess from PBH-star interactions leading to GRBs and MQs. For PBHs constituting DM, our signals will originate from locations where DM as well as NS/WD densities are the highest. Hence, we focus on the Galactic Center (GC). The fraction of dark matter in the form of PBHs that is required to explain $N_{\text{source}}$ positron sources (GRBs or MQs) necessary to address the excess can be expressed as

$$\left( \frac{\Omega_{\text{PBH}}}{\Omega_{\text{DM}}} \right) = \frac{N_{\text{source}}}{F_0 \tau_{\text{diff}} N_{\text{star}}} ,$$

where $F_0$ is the PBH-star capture rate (NS or WD) as before, $N_{\text{star}}$ is the respective compact star (NS or WD) population and $\tau_{\text{diff}} \simeq 3 \times 10^6$ yrs is the diffusion time for $\sim 100$ GeV positrons that sets the interaction timescale.

For MQs, contributions can come either from WDs or NSs that have been transformed into BHs. The amount of WD interacting binaries residing at the GC, which we take here as cataclysmic variables, is $N_{\text{WD}} \sim 10^5$ \cite{104}. The WD PBH capture rate is several orders below that of NS and we estimate that probability for PBH-WD MQs to account for the positron excess is low. NS capture rate as well as NS-binary GC population is higher and we find that NS-PBH MQs could provide a visible contribution to the excess. However, as we alluded to before, NS XRBs will already be emitting and contributing positrons even prior to PBH interactions. While the pre- and post-PBH capture jet emission from these systems is not expected to be identical, we shall not discuss this possibility further.

In order to explain the excess through PBH-induced GRBs, we require \cite{93} formation of $N_{\text{GRB}} = (8 \text{ kpc}/1 \text{ kpc})^2 = 64$ sources at the GC. For the GC milli-second pulsar population we take $N_{\text{NS}} \simeq 10^7$ stars \cite{27}. Since milli-second pulsars originate from binaries, we double the capture rate to approximate the increase due to stronger gravitational potential. For our fit to Eq. (5) we have scanned over a broad range of various astrophysical input parameters as follows. The DM density is varied as $50 \text{ GeV/cm}^3 < \rho_{\text{DM}} < 8.8 \times 10^2$ GeV/cm$^3$. Here, the lower bound corresponds to the “flat-core” Burkert profile \cite{105} with a uniform density in the central kpc region. The upper bound is the volume-averaged maximum allowed mass of the DM within 0.1 kpc of the GC, which is set by the criterion that DM does not exceed the baryonic content of $\sim 10^6 M_\odot$. The DM velocity dispersion values are considered to lie in the 50 km/s < $\bar{v}$ < 200 km/s range. The lower $\bar{v}$ limit corresponds to a possible DM disk within the halo \cite{106,107}, while the upper limit corresponds to the Navarro-Frenk-White DM density profile without adiabatic contraction \cite{108}. We further have included the effects of natal pulsar kicks, modifying the capture rate $F_0$ according to the method outlined in \cite{27}. We consider pulsar velocity dispersion between 48 km/s \cite{109} and 80 km/s \cite{110}. The final fit results are shown in Fig. 2 along with the current experimental constraints. We note that the considered timescales assume that a stellar BH will form very rapidly after PBH has been gravitationally captured. For $M_{\text{PBH}} \lesssim 10^{20}$ g the time for PBH to settle inside the star (post-capture) will exceed the interaction time set by $\tau_{\text{diff}}$. On the other hand, lower mass PBHs that have been captured earlier will then contribute at the same rate. Such behavior occurs until the settle time will reach the Galactic lifetime of $\sim 10^{10}$ yrs, which will happen for $M_{\text{PBH}} \lesssim 10^{18}$ g. At lower masses, however, the capture rate and settle time are not well understood and we thus show the full parameter range. As can be seen, if PBHs constitute a few percent or more of the DM they can significantly contribute to the positron excess through GRBs. We stress the effect of large uncertainties in the input quantities.

### Conclusions

We have proposed several scenarios how PBHs can lead to GRBs and MQs through interactions with compact stars. Further, we have argued that such interactions can generically lead to accretion disk formation, which is a major component of standard GRB production mechanisms. Relativistic outflows from GRBs as well as MQs will accelerate positrons, which will diffuse. We have found that if PBHs constitute a few percent of the DM or more, they can account or at least provide a significant contribution to the positron excess observed by multiple astronomy experiments. Our proposal combines positive features of astrophysical sources with dark matter, which have been typically disjoint in other proposals. The resulting parameter space for PBHs to address the excess is consistent with the parameter space where PBHs can contribute to the 511-keV GC line, $\tau$-process material abundance in Milky Way and dwarf spheroidal galaxies as well as fast radio bursts \cite{27}. PBHs can thus shed light on multiple astronomy puzzles simultaneously.

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Supplemental Material

ACCRETION DISK FORMATION

Accretion disk formation in a collapsing star system can be analyzed by estimating the amount of material residing outside the ISCO radius of the resulting black hole [61, 111]. Rotating milli-second pulsars spin near mass-shedding limit and can be described analytically by an extended Roche spheroid model [61, 111]. The equatorial radius $R_{\text{eq}}$ is stretched and exceeds the polar radius $R_p$, with $R_{\text{eq}} = (3/2)R_p$. Modeling the rapidly rotating star as a polytrope of index $n$, the spherical-coordinate density of the extended envelope is given by [61]

$$
\rho(r, \theta) = \frac{\xi_1^{3-n}(\xi_1^2|\theta'|(\xi_1))^{n-1} M}{4\pi R_p^3} \times \left(\frac{R_p}{r} - 1 + \frac{4}{27} \frac{r^2}{R_p^2} \sin^2(\theta)\right)^n ,
$$

where $\theta(\xi)$ is a solution to the Lane-Emden equation and $\xi$ denotes dimensionless radius, with $\xi_1$ corresponding to the star’s radius extent. The surface boundary is located along the $\rho = 0$ curve and can be described as [112]

$$
r(\theta) = \frac{3R_p \sin(\theta/3)}{\sin(\theta)} .
$$

We estimate the amount of angular momentum and mass located within the resulting BH that has consumed the star from within up to the polar radius $R_p$ through spherical accretion. At that point a wholesome NS is no longer present. With rotational symmetry in play, the ratio of the resulting BH mass to the original NS mass can be found from [111]

$$
\frac{M_{\text{BH}}}{M_{\text{NS}}} = \frac{\int_0^{\pi/2} \int_0^{R_p} \rho(r, \theta) r^2 \sin(\theta) d\theta d\theta}{\int_0^{\pi/2} \int_0^{r(\theta)} \rho(r, \theta) r^2 \sin(\theta) d\theta d\theta} .
$$

To make further progress, we need to assume a NS EoS or in our context a corresponding polytropic index. As discussed in the text, recent observations favor stiffer (lower $n$) NS profile. Since analytic solutions for the Lane-Emden equation are only known for $n = 0, 1, 5$, we take $n = 1$ as a NS description. We thus find that the mass ratio of the resulting BH compared to the original NS is $M_{\text{BH}}/M_{\text{NS}} \simeq 0.9$, which for $M_{\text{NS}} = 1.5M_\odot$ results in $M_{\text{BH}} \simeq 1.4M_\odot$.

As before, we take that all of the angular momentum from the in-falling particles is not transfered outside but is acquired by the BH instead. Hence, the resulting BH angular momentum is $J_{\text{BH}} = (2/5)\Omega_{\text{NS}}M_{\text{BH}}R_p^2$, yielding a spin parameter of $a \simeq 0.6$. The corresponding ISCO radius $r_{\text{ISCO}}(a = 0.6) \simeq 4GM_{\text{BH}}$ is found from [111]

$$
r_{\text{ISCO}} = GM_{\text{BH}}\left[3 + Z_2 - \sqrt{(3 - Z_1)(3 + Z_1 + 2Z_2)}\right],
$$

where

$$
Z_1 = 1 + (1 - a^2)^{1/3}(1 + a^{1/3} + (1 - a)^{1/3})
$$

and

$$
Z_2 = (3a^2 + Z_1^{1/2}).
$$

The final equatorial radius after BH has consumed the star up to $R_p$ is given by $R'_\text{eq} = R_{\text{eq}} - R_p + R_+$, where $R_+$ is the outer Kerr BH horizon as before. The ratio of the resulting disk mass $\Delta M$ to the original star’s mass $M_{\text{NS}}$ can then be found through summation of the remaining cylindrical mass shells outside the ISCO radius as proposed in [61]:

$$
\frac{\Delta M}{M_{\text{NS}}} = \xi_1^{3-n}(\xi_1^2|\theta'|(\xi_1))^{n-1} \times \int_{r'_{\text{ISCO}}}^{R'_{\text{eq}}} \int_0^{z'(r)} r\left(\frac{1}{\sqrt{r^2 + z^2}} - 1 + \frac{4}{27}r^2\right)^n dz dr ,
$$

where the prime on top of quantities in limits of integration denotes normalization to $R_p$ as $R'_{\text{eq}} = R_{\text{eq}}/R_p$, $r'_{\text{ISCO}} = r_{\text{ISCO}}/R_p$ and $z'(r) = z(r)/R_p$. Here, $z(r)$ describes the height of a cylindrical shell and can be in principle obtained by changing coordinates and then inverting Eq. (7). We approximate the height by performing the $dz$ integration up to $R'_+/2$, with $R'_+ = R_+/R_p$. We thus obtain $\Delta M/M_{\text{NS}} \simeq 0.1$, with the resulting disk mass of $M_{\text{disk}} \simeq 0.1M_\odot$. 

\[ \text{[61, 111]} \]
et al. (Virgo, LIGO Scientific), Phys. Rev. D87, 061101 (2013), arXiv:1301.4984 [astro-ph.CO].

[29] Y. B. Zeldovich and I. D. Novikov, Sov. Astron. 10, 602 (1967).

[30] S. Hawking, Mon. Not. Roy. Astron. Soc. 152, 75 (1971).

[31] B. J. Carr and S. W. Hawking, Mon. Not. Roy. Astron. Soc. 168, 399 (1974).

[32] J. Garcia-Bellido, A. D. Linde, and D. Wands, Phys. Rev. D54, 6040 (1996) arXiv:astro-ph/9605094 [astro-ph].

[33] M. Yu. Khlopov, Res. Astron. Astrophys. 10, 495 (2010) arXiv:0801.0116 [astro-ph].

[34] P. H. Frampton, M. Kawasaki, F. Takahashi, and T. T. Yanagida, JCAP 1004, 023 (2010) arXiv:1001.2308 [hep-ph].

[35] M. Kawasaki, A. Kusenko, and T. T. Yanagida, Phys. Lett. B711, 1 (2012) arXiv:1202.3848 [astro-ph.CO].

[36] M. Kawasaki, A. Kusenko, Y. Tada, and T. T. Yanagida, Phys. Rev. D94, 083523 (2016) arXiv:1606.07631 [astro-ph.CO].

[37] E. Cotner and A. Kusenko, (2016) arXiv:1612.02529 [astro-ph.CO].

[38] B. Carr, F. Kuhnel, and M. Sandstad, Phys. Rev. D94, 083504 (2016) arXiv:1607.06777 [astro-ph.CO].

[39] K. Inomata, M. Kawasaki, K. Mukaida, Y. Tada, and T. T. Yanagida, (2016) arXiv:1611.06130 [astro-ph.CO].

[40] K. Inomata, M. Kawasaki, K. Mukaida, Y. Tada, and T. T. Yanagida, (2017) arXiv:1701.02544 [astro-ph.CO].

[41] J. Georg and S. Watson, (2017) arXiv:1703.04825 [astro-ph.CO].

[42] J. L. Feng, Ann. Rev. Astron. Astrophys. 48, 495 (2010) arXiv:1003.0994 [astro-ph.CO].

[43] G. Bertone, D. Hooper, and J. Silk, Phys. Rept. 405, 270 (2005) arXiv:hep-ph/0404175 [hep-ph].

[44] S. Bird et al., Phys. Rev. Lett. 116, 201301 (2016) arXiv:1603.00464 [astro-ph.CO].

[45] T. Nakamura, M. Sasaki, T. Tanaka, and K. S. Thorne, Astrophys. J. 487, L139 (1997) arXiv:astro-ph/9708060 [astro-ph].

[46] S. Clesse and J. Garcia-Bellido, Phys. Rev. D92, 023524 (2015) arXiv:1501.07565 [astro-ph.CO].

[47] M. Raidal, V. Vaskonen, and H. Veermäe, (2017) arXiv:1707.01480 [astro-ph.CO].

[48] Yu. N. Eroshenko, (2017) arXiv:1604.04932 [astro-ph.CO].

[49] M. Sasaki, T. Suyama, T. Tanaka, and S. Yokoyama, Phys. Rev. Lett. 117, 061101 (2016) arXiv:1603.08338 [astro-ph.CO].

[50] S. Clesse and J. Garcia-Bellido, (2016) arXiv:1610.08479 [astro-ph.CO].

[51] V. Tikhovtsov, (2017) arXiv:1707.05849 [astro-ph.CO].

[52] B. P. Abbott et al. (Virgo, LIGO Scientific), Phys. Rev. Lett. 116, 061102 (2016) arXiv:1602.03837 [gr-qc].

[53] B. P. Abbott et al. (Virgo, LIGO Scientific), Phys. Rev. Lett. 116, 241103 (2016) arXiv:1606.04855 [gr-qc].

[54] B. P. Abbott et al. (VIRGO, LIGO Scientific), Phys. Rev. Lett. 118, 221101 (2017) arXiv:1706.01812 [gr-qc].

[55] G. M. Fuller, A. Kusenko, and V. Tikhovtsov, Phys.
