Weakly interacting massive particles and neutron stars

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Neutron stars are used to set constraints on the characteristics of weakly interacting massive particles (WIMP's) suggested as dark-matter candidates. Some special classes of WIMP's are ruled out because they would be trapped in neutron stars, concentrate towards the star center, and become self-gravitating. This results in the formation of a mini black hole that consumes the neutron star, transforming it into a black hole, on a time scale shorter than observed ages of neutron stars in various astrophysical systems.

I. INTRODUCTION

Many astrophysical observations indicate "missing mass," or "missing light," on galactic as well as larger scales. The observed luminous matter, plus the accounted-for nonluminous matter, is not enough to explain the dynamics on these scales.[1] An attractive and widely discussed possible explanation is that this mass is in the form of an ambient gas of relic weakly interacting massive particles (WIMP's). WIMP's with masses $m_x \gtrsim 10$ GeV are of particular interest for several reasons.

(i) They constitute "cold dark matter" which, unlike "hot dark matter" (e.g., light, ≤ 50 eV, neutrinos), readily clusters on galactic scales.[2,3]

(ii) If $X^0 \neq \bar{X}^0$ and the X^0 and $\bar{X}^0$ particles annihilate in the early Universe leaving a small fractional excess $\delta_X \equiv (n_X - n_{\bar{X}})/(n_X + n_{\bar{X}})$, then the ratio of the baryonic to the dark-matter density is

$$\frac{\rho_B}{\rho_X} \sim \frac{\delta_B}{\delta_X} \frac{m_B}{m_X} \frac{n_B + n_{\bar{B}}}{n_X + n_{\bar{X}}} , \tag{1}$$

where δ_B is the baryonic fractional excess. The observed value $\rho_B/\rho_X \sim 0.01-0.1$ follows naturally if $m_X \sim (10-10^3)m_B$, $n_{X+\bar{X}} \sim (1-0.1)n_{B+\bar{B}}$, and $\delta_X \sim (1-0.1)\delta_B$. There is no convincing explanation for the baryon asymmetry, $\delta_B \sim 10^{-9}$, but it is reasonable to suppose that a similar mechanism generates δ_X and hence $\delta_X \sim \delta_B$. We will address subsequently the conditions required for sufficient $X\bar{X}$ annihilation, so that the mass density in $X\bar{X}$ does not exceed the cosmological closure density.

(iii) $m \sim 10^2$ GeV is the scale of $SU(2) \times U(1)$ symmetry breaking and of the conjectured supersymmetry breaking.

(iv) WIMP's of mass m_X moving at a galactic virial velocity, $\sim 10^{-3}c$, deposit in a nuclear collision, a recoil energy of the order of

$$\Delta E \sim \min(m_X, m_{A,Z}) \frac{v_{\rm vir}^2}{2} \sim 10 \text{ keV} \tag{2}$$

for m_X, $m_{A,Z} \geq 20$ GeV. This exceeds the detection threshold for existing germanium ($A = 80$) devices and planned granular superconducting devices.[4,5] The large cross section for ordinary weak (Z^0 exchange) interaction with nuclei[4] (for $A \sim 10$),

$$\sigma_{Z^0} \sim \sin^2\theta_W \frac{G_F^2}{\pi\hbar^4} \frac{(m_X m_{A,Z})^2}{(m_X + m_{A,Z})^2} (A-Z)^2 \sim 10^{-35} \text{ cm}^2 , \tag{3}$$

allows excluding, at present, WIMP's of $m_X \geq 20$ GeV and $\sigma_{X\,\rm nuclear} \sim \sigma_{Z^0}$. WIMP's of smaller mass and/or cross sections could be accessible in future, more sensitive experiments. Strictly speaking, these experiments do not exclude very massive WIMP's, say, $m_X > 10^6$ GeV. Since the WIMP mass density is fixed to be equal to the mass density of the dark matter, the WIMP flux scales as m_X^{-1}, so that for very massive WIMP's the fluxes would be too low to be detected.

(v) WIMP's with mass $\gtrsim 4$ GeV and nuclear cross sections somewhat larger than σ_{Z^0}, could accumulate in the Sun, lower the temperature in the Sun center so as to affect the flux of energetic neutrinos, and resolve the solar-neutrino problem.[6–8] In certain circumstances the accumulated WIMP's could also affect the life span of the host star.[9,10]

In this paper we examine whether neutron stars can be used to rule out WIMP's with cross sections that are smaller than those accessible to the experiments. We find that only WIMP's of very special characteristics can be ruled out. We show that massive bosonic WIMP's ($m_X \gtrsim 200$ GeV), without coherent vectorial Z^0-exchange interactions, can coalesce at the neutron-star core and form a mini black hole. If one takes into account also the possibility of a Bose condensation, the limit can be reduced to $m_X \gtrsim 1$ GeV. The mini black hole will consume the host neutron star and transform it into a black hole. The overall time scale is shorter than observed ages of neutron stars in various astrophysical systems. Therefore, these WIMP's are ruled out. Very massive fermionic WIMP's ($m_X \gtrsim 10^7$ GeV) could also form a mini black hole on Hubble time scale. However such massive WIMP's are ruled out by the requirement that the mass density of the symmetric component $\min\{n_X, n_{\bar{X}}\}$, left over from the annihilation in the early Universe, be smaller than the cosmological critical density. However, WIMP's with strong coherent scalar-exchange attractive interaction, that overcomes the Fermi degeneracy repulsion and the repulsive Z^0-exchange interaction, are also ruled out.

The envisioned scenario above, proceeds via the following stages.

(A) Accumulation stage: The neutron star accumulates the WIMP's which thermalize and assume a Maxwellian distribution, under the influence of the gravitational field of the neutron star. As the star cools, the thermal radius of the WIMP distribution decreases while the total mass within it increases, until at a time t_A they become self-gravitating.

(B) (Mini-) black-hole formation: The self-gravitating WIMP sphere forms a mini black hole on a time scale t_B.

(C) Consumption of the neutron star: The mini black hole accretes matter from the dense core of the neutron star and grows. The explosive growth process winds up at t_C in a complete transformation of the host neutron star into a black hole.

We investigate under which conditions will the above scenario occur, and use the existence of neutron stars to rule out those WIMP's for which it does. Clearly, for the above scenario to happen, t_A, t_B, and t_C should be shorter than the Hubble time $t_0 = H_0^{-1} \sim 1-2\times 10^{10}$ yr. These time spans depend on the parameters of the neutron star as well as on the WIMP's density and velocity and on the cross sections for various processes involving WIMP's.

The relevant parameters of the neutron star are the radius R, the stellar mass M, the core density ρ_c, and the average stellar density $\bar{\rho} = [3/(4\pi)]MR^{-3}$. We will use as representative values[11] $\rho_c = 10^{15}$ g cm^{-3}, $R = 10^6$ cm, $M = 1.4 M_\odot$ implying $\bar{\rho} = 6.7\times 10^{14}$ g cm^{-3}.

The dark-matter density in the solar neighborhood is estimated to be[1,12] $\sim 0.11 M_\odot$ pc$^{-3} \sim 7.4\times 10^{-24}$ g cm^{-3}. This however is the density of the disk component of the galactic dark matter. Since the WIMP's, due to their small cross sections, do not dissipate their energy, they cannot be confined to the galactic disk and are candidates for the halo component of the dark matter. Thus, for the solar neighborhood we take[7] $\rho_x \sim 0.01 M_\odot$ pc$^{-3} \sim 6.7\times 10^{-25}$ g cm^{-3}. For the WIMP velocity, relative to the neutron star, we adopt a value of the order of the solar neighborhood rotation velocity, thus $v_X \sim 250$ km sec^{-1}. Assuming a spheroidal dark-matter density distribution with $\rho \propto r^{-2}$ within R_0, the solar distance from the galactic center (8.5–10 kpc), we can estimate that at a distance of $\frac{1}{2}R_0$ the corresponding value of the density is ~ 4 times larger while the velocity is roughly the same as before. In case that the WIMP spatial distribution is clumped, one expects regions with even higher density. If the clumping is on too small scales it will, on the average, decrease the flux captured by a given neutron star. Such primordial clumping on subgalactic scales is quite possible due to the small Jeans mass associated with such heavy particles.

The X-nucleon cross section will be parametrized in terms of an effective Fermi constant G_F' by scaling down from the standard weak cross section:

$$\sigma_{Xn} \sim \frac{1}{\pi\hbar^4} m_N^2 G_F'^2 \sim \epsilon 10^{-38}\ \text{cm}^2 \qquad (4)$$

with $\epsilon = G_F'^2/G_F^2 < 1$ and $m_N \sim 1$ GeV being the nucleon mass. The assumption here is that the interaction is generated by a heavy t-channel exchange of some particle Y with $m_Y \gg m_N$. Note that for $m_X \gg m_N$ the cross section is independent of m_X.

II. X-$\bar{X}$ ANNIHILATION IN THE EARLY UNIVERSE

Our program is to use the existence of neutron stars in order to rule out those WIMP's for which the general scenario, outlined above, takes place. First we must make sure that annihilation in the early Universe was efficient enough so that the symmetric WIMP number density $n_X^s = \min\{n_X, n_{\bar{X}}\}$ does not contribute a mass density exceeding the cosmological critical density. Note that only the excess $n_X - n_{\bar{X}}$ is relevant to the scenario described above since the symmetric component will annihilate in the neutron star.

Annihilation becomes relevant when the cosmological temperature satisfies $kT \sim m_X c^2$ so that the back reactions "radiation"$\rightarrow X\bar{X}$ become exponentially damped. So long as the annihilation rate exceeds the rate of the cosmological expansion, the WIMP number density is given by the equilibrium expression

$$\frac{n_X^s}{n_\gamma} \sim \frac{g}{4}(m_X c^2/kT)^{3/2} e^{-(m_X c^2/kT)} , \qquad (5)$$

where g is the number of spin states of the X particles, and n_γ, the photon number density, is given by

$$n_\gamma = 0.244(kT/\hbar c)^3 . \qquad (6)$$

The annihilation effectively stops, and the number density of the symmetric component becomes frozen-in, when the rate of the annihilation becomes smaller than the rate of the cosmological expansion. The annihilation rate is

$$\tau_a^{-1} = n_X^s \sigma_a v_X \sim g_a^4 n_X^s \hbar^2 m_X^{-2} c^{-1} , \qquad (7)$$

where the annihilation cross section σ_a is parametrized by $\sigma_a = (v_X/c)^{-1} g_a^4 [\hbar/(m_X c)]^2$. Such a parametrization is indeed natural. The annihilation must involve an exchange of a particle X' carrying the X quantum numbers but heavier than the X itself. Otherwise, the decay $X \rightarrow X' +$light particles would be possible and the X particles could not have survived over a Hubble time scale. In case that the X couple to Z^0, one has $X\bar{X} \rightarrow$virtual $Z^0 \rightarrow$light particles, with the same cross section, σ_a given above.

The rate of the cosmological expansion is given by[13]

$$R^{-1} dR/dt = \left[\frac{8\pi}{3}\frac{G}{c^2}\rho\right]^{1/2} = \left[\frac{4\pi}{3} N \frac{G}{c^2} a T^4\right]^{1/2} , \qquad (8)$$

where $a = (\pi^2/15)k^4\hbar^{-3}c^{-3}$ is the Stephan-Boltzmann constant and N is the number of degrees of freedom of all the particles that are relativistic at that epoch. By equating the above two rates we obtain

$$\frac{n_X^s}{n_\gamma}=5.57\times10^{-19}\, g_a^{-4}N^{1/2}(m_Xc^2/kT)(m_X/1\text{ GeV}) . \tag{9}$$

By equating n_X^s/n_γ of (9) with that of (5) one gets a relation between the particle mass and the freeze-in temperature:

$$\frac{g}{4}(m_Xc^2/kT)^{3/2}e^{-(m_Xc^2/kT)} =5.57\times10^{-19}g_a^{-4}N^{1/2}(m_Xc^2/kT)(m_X/1\text{ GeV}) . \tag{10}$$

We impose now the condition that the present epoch mass density in WIMP's does not exceed the critical cosmological density. Thus for $H_0 =100$ km sec^{-1}Mpc^{-1}, $(n_X^s/n_\gamma)_{\text{now}}\lesssim 1.4\times10^{-8}(m_X/1\text{ GeV})^{-1}$. Note also that the present epoch photon-number density is larger by a factor $\sim N$, due to annihilations that took place between the annihilation epoch of the X particles and the present epoch. Using these in (9) yields

$$m_X \lesssim 1.6\times10^5\text{ GeV}g_a^2N^{1/4}\left[\frac{kT}{m_Xc^2}\right]^{1/2} . \tag{11}$$

Taking $N\sim100$ we get from (10) and (11) that $m_Xc^2/kT\gtrsim31$, which when used in (11) results in

$$m_X \lesssim 0.95\times10^5 g_a^2\text{ GeV} . \tag{12}$$

III. WIMP'S ACCRETION ONTO THE NEUTRON STAR

To find the X-particle cross section for collision with the neutron star, consider a particle at infinity with velocity v_X and impact parameter b, whose closest approach to the star is R; the sought-for cross section is πb^2. Using the equations of test particles trajectories in the Schwarzschild metric[14] we obtain, for a particle just grazing the star surface,

$$1=\frac{(E_\infty/c^2)^2}{1-2GM/(c^2R)}-\frac{L^2}{R^2c^2} , \tag{13}$$

where $L=v_Xb$ and $E_\infty=(1-v_X^2/c^2)^{-1/2}c^2$ are the conserved angular momentum and the conserved energy (as measured by an observer at infinity) per unit mass, respectively. From (13) follows (using $v_X^2/c^2\ll1$)

$$b=(2GMRv_X^{-2})^{1/2}(1-2GM/c^2R)^{-1/2} =R\frac{v_{\text{es}}}{v_X}(1-2GM/c^2R)^{-1/2} , \tag{14}$$

where $v_{\text{es}}=(2GM/R)^{1/2}$ is the escape velocity as measured by an observer at rest on the neutron-star surface. From (14) we find the mass flux of WIMP's colliding with the star

$$\frac{dM_X}{dt}=\pi b^2\rho_Xv_X=2\pi RGMv_X^{-1}\rho_X(1-2GM/c^2R)^{-1}. \tag{15}$$

Using the previously adopted solar neighborhood values of ρ_X and v_X, we get a flux of ~53 g sec^{-1}. At half a distance from the galactic center the flux is estimated to be ~4 times larger.

Let us find the conditions required for the WIMP's to lose enough energy while traversing the neutron star, so as to become bound. For $m_X\gg m_N$, WIMP's lose in an average nuclear collision a fraction $2m_N/m_X$ of their kinetic energy, as measured by an observer at rest at the location of the collision. Because the WIMP's velocity at the neutron star is so much larger than v_X, even one collision can lead to sufficient energy loss and bound the WIMP. In order that a passage through the star will bound the WIMP it is required that $n_{\text{col}}\geq1$ and that the total energy (including rest mass), as viewed from infinity, be smaller than m_Xc^2. Note, that if the latter condition is met on the first passage through the star, a bound orbit results. In subsequent revolutions the particle must pass through the star, regardless of the orbit periastron precessions. Taking into account the redshift correction, the condition for binding is found to be

$$2n_{\text{col}}(m_N/m_X)m_Xc^2[1-(1-2GM/c^2R)^{1/2}]\gtrsim\tfrac{1}{2}m_Xv_X^2 . \tag{16}$$

Note that (15) and (16) are the general-relativistic generalizations of the classical Newtonian results.

The number of X-nucleon collisions, for a typical path through the star is

$$n_{\text{col}}\sim Rn\sigma_{Xn} . \tag{17}$$

For the cross section of Eq. (4) and the parameters adopted in Sec. I, we find that $n_{\text{col}}\geq1$ and the condition for binding, Eq. (16), are satisfied if, respectively,

$$\epsilon\geq2.5\times10^{-7},\quad m_X\leq5.2\times10^{12}\epsilon\text{ GeV} . \tag{18}$$

Because the WIMP velocity at the neutron star is semi-relativistic, there is no suppression of the cross section due to neutron degeneracy. The latter will be relevant after the trapped WIMP's thermalize and become nonrelativistic.

The conditions (18) are satisfied for a wide range of m_X and ϵ (<1). Thus, in order to capture WIMP's in a neutron star one does not need a "stronger than weak" cross section that is required to capture efficiently "light," few GeV, WIMP's in a main-sequence star.[6–10] Note that for a constant star mass, the critical mass in the second equation of Eq. (18) is proportional to R^{-3}, so that capture in a white dwarf is limited to WIMP's of lower masses. Capture by a main-sequence star of particles more massive than few GeV, requires that $\epsilon>1$. If the first of the conditions of (18) is not satisfied, then only a fraction, $0.4\epsilon10^7$, of the incident WIMP's will suffer at least one collision. Henceforth we will assume that both conditions of (18) are satisfied so that (15) indeed expresses the actual rate of mass accretion onto the neutron star. The mass accreted during a time t is thus

$$M_X(t)=1.6\times10^{16}\text{ g}(t/10^7\text{ yr}) . \tag{19}$$

Note that for neutron stars at a distance from the galactic center equal to half that of the solar system, the accreted mass is a factor ~ 4 larger.

IV. THERMALIZATION OF THE WIMP'S IN THE NEUTRON STAR AND ONSET OF SELF-GRAVITY

From numerical solutions of neutron-star models[15] follows that the density in the core is almost constant up to a radius of $\sim 2-4$ km, depending on the nuclear matter equation of state. We will be interested in the gravitational field up to a radius of ~ 1 m so that we take the density to be constant, ρ_c. We find below, that even though the star center is in a very deep gravitational potential well, the gravitational field, in the vicinity of the star center, can be treated as a Newtonian field. Introducing a metric in spherical-polar coordinates

$$ds^2 = e^{2\phi}dt^2 - e^{2\lambda}dr^2 - r^2(d\theta^2 + \sin^2\theta\, d\Phi^2) \tag{20}$$

the solutions of the Einstein field equations[14] near $r=0$, in the constant density region, are

$$e^{2\lambda} = \left[1 - \frac{8\pi G\rho_c}{3c^2}r^2\right]^{-1}, \qquad \phi = \phi_c + \frac{2\pi G}{3c^2}(\rho_c + 3p_c)r^2 . \tag{21}$$

By rescaling the time coordinate, $d\bar{t} = e^{\phi_c}dt$, one gets a metric representing a weak Newtonian gravitational field of a potential $\bar{\phi} = \phi - \phi_c$. Note that the pressure contributes to the field and (depending on the equation of state) $3p_c \sim \rho_c$. For the adopted value of $\rho_c = 10^{15}$ g cm^{-3},

$$\bar{\phi} \sim \frac{4\pi G\rho_c}{3c^2}r^2 \sim 3\times 10^{-9}(r/1\text{ m})^2 . \tag{22}$$

Thus, for $r<1$ m, indeed $\bar{\phi} \ll 1$ and the gravitational field can be treated as Newtonian. Note however, that the star center is at the bottom of a very deep gravitational potential well. This should be taken into account when transforming various quantities from the frame of a local observer at the neutron-star center to that of an observer at infinity, by employing the redshift factor $z_c = e^{-\phi_c} - 1$ (typically, the latter ~ 1).

Initially, the self-gravity of the WIMP's is negligible compared to that of the neutron star and Eq. (22) correctly represents the gravitational field acting on the WIMP's. The WIMP's collide with the neutrons and as a result they attain the neutron matter temperature. Let us assume a thermal equilibrium via these collisions and estimate the time scale for thermalization in order to check the consistency of this assumption. Since the neutrons are degenerate, only those neutrons with momentum $\mathbf{p}$ satisfying $|\mathbf{p}+\Delta\mathbf{p}| > p_F$, with $\Delta\mathbf{p}$ the momentum gained in the collision, can participate in the collision. This results in a reduction of the cross section by a factor $3\Delta p/4p_F$. Since the X particles are far more massive than the neutrons, a collision of an X particle of velocity v_X, the assumed thermal velocity, with a neutron will typically impart that velocity to the latter. Therefore, $\Delta p \sim m_N v_X$, with $v_X \sim (3kT/m_X)^{1/2}$. In each collision, on the average, a fraction $2m_N/m_X$ of the energy is lost by the X particle. The time scale, for losing an energy of the order of the initial energy, is therefore ($p_F c \sim 520$ MeV for $\rho_c = 10^{15}$ g cm^{-3})

$$\tau_{\text{col}} \sim \frac{2m_X}{9m_N}\frac{m_X c^2}{kT}\frac{p_F}{m_N c}(n\sigma_X c)^{-1} \sim 0.22\text{ yr}(\epsilon/10^{-5})^{-1}(T/10^5\text{ K})^{-1}(m_X/1000\text{ GeV})^2 . \tag{23}$$

Thus, for a very wide range of m_X and ϵ, the time scale for thermal relaxation is very short, compared to the $\sim 10^7$ yr it takes for the interior temperature T to get down to 10^5 K (Refs. 16 and 17). We may therefore consider the WIMP's to be in thermal equilibrium at the temperature of the interior of the neutron star. The latter is isothermal[16,17] due to the large thermal conductivity of the degenerate neutrons.

So long as the WIMP average separation, $n_X^{-1/3}$ is larger than the thermal de Broglie wavelength,

$$\lambda = \hbar/(kTm_X)^{1/2} \sim 10^{-11}\text{ cm}(T/10^5\text{ K})^{-1/2}(m_X/1000\text{ GeV})^{-1/2} ,$$

the WIMP's obey the classical Maxwell Boltzmann statistics. As discussed below, for bosonic WIMP's a possible Bose condensation can assist the onset of self-gravity while for fermionic WIMP's, the collapse of the self-gravitating WIMP's is strongly impeded unless the WIMP's are very massive. For $n_X^{-1/3} \geq \lambda$, the WIMP distribution in the harmonic gravitational potential of the neutron star is Maxwellian both in velocities and in distances from the center. Therefore, the thermal rms value of the WIMP distance from the center is

$$r_{\text{th}} \sim \left[\frac{9KT}{8\pi G\rho_c m_X}\right]^{1/2} \sim 6.45\text{ cm}(T/10^5\text{ K})^{1/2}(m_X/1000\text{ GeV})^{-1/2} . \tag{24}$$

Let us next estimate the time required to accumulate a sufficiently large total mass in WIMP's, M_X, so that it will become self-gravitating. By definition, this is the time scale t_A mentioned before. The WIMP's will become self-gravitating when their total mass within r_{th} exceeds the ambient neutron star mass within the same radius. Using (19) and (24) we get

$$t_A = \frac{4\pi}{3}(dM_X/dt)^{-1}\rho_c r_{\text{th}}^3 \sim 7.3\times 10^8\text{ yr}(T/10^5\text{ K})^{3/2}(m_X/1000\text{ GeV})^{-3/2} . \tag{25}$$

Note that for the range of m_X allowed by the considerations of annihilation in the early Universe, Eq. (12), the time scale for thermalization, $\tau_{\rm col}$ of (23), is shorter than t_A so that the latter determines the time when the WIMP's will become self-gravitating. Another estimate, based on a detailed model for a stationary Boltzmann WIMP density distribution in the combined gravitational field of the neutron star and the WIMP's (A. Gould, private communication), yields that the maximal WIMP mass which is stable against self-collapse is $\sim\frac{4}{3}$ larger than what we get here.

If the X particles are bosons, self-gravity can set in earlier than t_A due to a Bose-Einstein condensation in the external gravitational potential of the neutron star. The condition for the occurrence of the condensation is[18] that $T<T_{\rm cr}$,

$$T_{\rm cr}=\frac{3.31\hbar^2 n_X^{2/3}}{g^{2/3}km_X} \tag{26}$$

with g denoting the number of spin states and $n_X=(M_X/m_X)/[(4\pi/3)r_{\rm th}^3]$ is the number density of the X particles. From (24) and (26) follows that for $g=1$, $T<T_{\rm cr}$ if the temperature T satisfies

$$T\lesssim 2.5\times10^5\ {\rm K}(t/10^7\ {\rm yr})^{1/3}(m_X/1000\ {\rm GeV})^{-1/3}\ . \tag{27}$$

The radius corresponding to the condensed ground state is

$$r_{\rm cond}=(\hbar/m_X)^{1/2}\left[\frac{32\pi}{3}G\rho_c\right]^{-1/4}$$
$$\sim 3.6\times10^{-6}\ {\rm cm}(m_X/1000\ {\rm GeV})^{-1/2}\ .$$

For the range of WIMP masses, $m_X\gtrsim 1$ GeV considered in the present work (the X-nucleon cross sections used apply for this range), this radius is small enough for the WIMP's, included within it, to be self-gravitating. From cooling curves of neutron stars follows that a time of $\sim10^7$ yr is required for the neutron star to cool to $\sim10^5$ K (Refs. 16 and 17). A comparison of (27) and (25) shows that for masses 1 GeV $\lesssim m_X\lesssim 1.7\times10^4$ GeV the condensation can indeed shorten the time required for the onset of self-gravity, compared to t_A of (25). This is so since the lower the value of m_X the larger are n_X and $T_{\rm cr}$. Thus, for the above-mentioned lower part of the WIMP mass range, Bose condensation determines the time scale for onset of self-gravity.

V. CONDITIONS FOR THE FORMATION OF A (MINI) BLACK HOLE

For $t\geq t_A$, the WIMP self-gravity dominates and the WIMP sphere is unstable against gravitational collapse. As we shall see in Sec. VI, the released gravitational energy is efficiently thermalized and transferred outward to the surrounding neutron star. In the absence of some strong impediment, the collapse will continue, forming eventually a mini black hole of mass $M_X\sim t_A dM_X/dt$.

In listing such impediments let us consider first the quantum kinetic energy arising from the confinement of the WIMP's to a sphere of radius r_X. The kinetic energy of a particle $E_k=(m_X^2c^4+p_X^2c^2)^{1/2}-m_Xc^2$ satisfies the inequality $E_k\leq p_Xc$. For a gravitational collapse to a black hole to happen, the gravitational energy per particle should be larger than the kinetic energy per particle: namely,

$$GM_Xm_X/r_X\gtrsim p_Xc\ . \tag{28}$$

For bosons $p_X\sim\hbar/r_X$ and (28) yields

$$m_X\gtrsim\hbar cG^{-1}M_X^{-1}\sim10^{-2}\ {\rm GeV}(t_A/10^7\ {\rm yr})^{-1} \tag{29}$$

which is satisfied by many orders of magnitude, for the range of WIMP masses considered here.

In the case of fermions $p_X\sim(M_X/m_X)^{1/3}\hbar/r_X$ and the condition for black-hole formation, (28), yields

$$m_X\gtrsim(\hbar c/G)^{3/4}M_X^{-1/2}$$
$$\sim4.3\times10^8\ {\rm GeV}(t_A/10^7\ {\rm yr})^{-1/2}\ . \tag{30}$$

Thus only very massive WIMP's, $m_X\gtrsim10^8$ GeV, could collapse to form a black hole in 10^8 yr. Even if t_A is allowed to be as large as the Hubble time only WIMP's of mass $\gtrsim10^7$ GeV could collapse to form a black hole. The requirements on annihilation in the early Universe, Sec. II, rule out such massive WIMP's. Therefore, fermions cannot form a mini black hole. Also composite bosonic WIMP's made of an even number of fermionic constituents, with mass smaller than the above masses, will not collapse to a black hole, unless the compositeness scale is smaller than

$$GM_Xc^{-2}N^{-1/3}\sim5.7\times10^{-25}\ {\rm cm}(t_A/10^7\ {\rm yr})^{2/3}$$
$$\times(m_X/1000\ {\rm GeV})^{1/3}\ .$$

Otherwise, the composite bosons will dissociate into their fermionic constituents before the WIMP sphere enters inside its horizon and becomes a black hole.

A repulsive coherent interaction

$$\hbar cg_W^2\sin^2\theta_W e^{-m_Zc\hbar^{-1}r}r^{-1}$$

between the X^0 particles, coupling to the Z^0 exchange, will stop the contraction once

$$r_X\lesssim g_W\sin\theta_W(\hbar/m_Xc)(m_{\rm Pl}/m_{Z^0})$$
$$\sim0.1\ {\rm cm}(m_X/1000\ {\rm GeV})^{-1}\ ,$$

where $m_{\rm Pl}\sim10^{19}$ GeV is the Planck mass. So long as this r_X is larger than the Schwarzschild radius of the WIMP sphere, $2GM_X/c^2\sim2.4\times10^{-12}$ cm$(t_A/10^7$ yr) the collapse to a black hole will be prevented. Even a vectorial coupling to a much more massive Z'^0 (up to 10^{13} GeV) could still prevent the collapse. A more careful examination reveals that the Z^0 repulsion could be overcome once the masses are large enough to overcome the Fermi barrier, see Eq. (30). The reason for this is that the "Z^0-charged" WIMP sphere can be readily neutralized by particles carrying the opposite Z^0-charge (e.g., neutrinos or neutrons). Since the number of these neutralizing particles equals that of the WIMP's, they will exert a Fermi pressure equal to the one that fermionic WIMP's would. Thus, heavy WIMP's (bosons or fermions but now with ordinary weak coupling, $\epsilon \sim 1$) that satisfy (30) will overcome this Fermi pressure and collapse to form a mini black hole. However, we note again that such massive WIMP's are ruled out by the requirements on the WIMP annihilation in the early Universe, Sec. II.

Coming back to the question of the annihilation, one may wonder whether the Z^0 attraction between $X\bar{X}$ can enhance their annihilation in the early Universe so that more massive WIMP's will be allowed. This attraction can lead to formation of a "Coulombic" bound state once $kT \lesssim g_W^2 \sin^2\theta_W m_{Z^0} c^2$. This requires either a three-body collision or the emission of an extra $\nu\bar{\nu}$, $e\bar{e}$, etc., pair. Once the bound state is formed, annihilation ($X\bar{X} \to Z^0 \to \nu\bar{\nu}, \ldots$) is guaranteed. However, the $X\bar{X}$ capture cross section, unless assisted by some other interactions, cannot enhance sufficiently the annihilation and the mass limit obtained in Sec. II is unchanged.

Finally, it is important to note that the Fermi and Z^0 repulsions could be overcome by an attractive force due to an exchange of a scalar of mass m_S and coupling g_S provided that $g_S^2/m_S^2 \geq g_W^2 \sin^2\theta_W / m_{Z^0}^2$.

VI. COOLING OF THE COLLAPSING WIMP SPHERE

In the absence of the above specific impediments to the collapse, let us consider the time scales associated with the collapse process. Consider a contraction of the radius of the WIMP sphere, r_X, to $r_X/2$. Such a contraction releases a gravitational energy $\Delta E \sim GM_X^2/r_X$. As shown next, the time scale for thermalization of the contracting self-gravitating WIMP sphere is shorter than the dynamical free-fall time scale τ_d:

$$\tau_d \sim (r_X^3/GM_X)^{1/2} \sim 6\times 10^{-5}\ \sec(r_X/r_0)^{3/2}\ , \quad (31)$$

where r_0 is the initial value of r_X, at $t=t_A$, when self-gravity sets in

$$r_0 = \left[t_A \frac{dM_X}{dt} \frac{3}{4\pi\rho_c} \right]^{1/3} \sim 1.6\ \mathrm{cm}(t_A/10^7\ \mathrm{yr})^{1/3}\ . \quad (32)$$

Therefore, the released gravitational energy is quickly thermalized and the collapse time scale will depend on the cooling time scale.

Let us consider the cross section for XX scattering which, as seen below, dominates the thermalization time scale of the WIMP's. For the sake of simplicity let us assume that the mass of the exchanged Y particle (see end of Sec. I) $m_Y \gg m_X$. In this case the XX cross section is

$$\sigma_{XX} \sim \epsilon \pi^{-1} \hbar^{-4} G_F^2 m_X^2$$
$$\sim \epsilon 10^{-32}\ \mathrm{cm}^2 (m_X/1000\ \mathrm{GeV})^2\ . \quad (33)$$

The time scale for thermalization is

$$\tau_{XX} = (n_X \sigma_{XX} v_X)^{-1} \quad (34)$$

yielding

$$\tau_{XX} \sim 6.4\times 10^{-4}\ \sec(\epsilon/10^{-5})^{-1}(m_X/1000\ \mathrm{GeV})^{-1}(t_A/10^7\ \mathrm{yr})^{-1/3}(r_X/r_0)^{7/2}\ , \quad (35)$$

where v_X is given by

$$v_X \sim (GM_X/r_X)^{1/2} \sim 2.6\times 10^4\ \mathrm{cm\,sec}^{-1}(t_A/10^7\ \mathrm{yr})^{1/3}(r_0/r_X)^{1/2}\ . \quad (36)$$

If $m_Y < m_X$ the cross section increases and τ_{XX} becomes even shorter. In either case, τ_{XX} is very short. For comparison, the time scale for thermalization via xn collisions is given by

$$\tau_{Xn} = \frac{m_X}{2m_N} \frac{1}{n\sigma_{Xn} v_X} \frac{4p_F}{3m_N v_X} \sim 1.3\times 10^8\ \sec(\epsilon/10^{-5})(m_X/1000\ \mathrm{GeV})(t_A/10^7\ \mathrm{yr})^{-2/3} \frac{r_X}{r_0}\ , \quad (37)$$

where the factor $m_X/2m_N$ is the average number of scatterings required for an X particle to transfer an energy of the order of its kinetic energy. The Fermi blocking factor for the neutrons further reduces the cross section.

Since $\tau_{XX} \ll \tau_{Xn}$, the thermalization time scale is τ_{XX}. From (31) and (35) follows that for $r_X/r_0 \lesssim 0.3$, the thermalization time scale is shorter than the dynamical time scale τ_d. Therefore, the collapse time will depend on the cooling time scales.

We consider, below, three cooling mechanisms for the collapsing WIMP sphere. The first is by X-neutron collisions. The neutrons transfer the energy to the surrounding neutron star via the large thermal conductivity of the degenerate neutrons. This mechanism requires no additional assumptions regarding the WIMP interactions, and as seen from Eq. (37), implies a cooling time scale which is very short compared to t_A and thus is sufficient by itself. We wish, however, to consider also the possibility that the WIMP's interact not only with neutrons but also with neutrinos and other superweakly interacting (e.g., axions) massless particles and can emit these particles as a blackbody radiation. The second cooling mechanism is thus operating when $r_X > l_X$, where l_X is the mean free path for XX collisions. In this case the surface of the WIMP sphere radiates as a blackbody. In this case, too, the heat is further transferred outward by the neutron thermal conductivity. The third cooling mechanism is neutrino cooling and is operative if the WIMP sphere is opaque to νX scattering so that a population of thermal $\nu\bar{\nu}$ is generated.

The time scale for the first mechanism is τ_{Xn} given by (37). The time scale for the second mechanism is estimated next. The luminosity from the surface of the WIMP sphere is

$$L_X = 4\pi\sigma r_X^2 T_X^4 \ , \tag{38}$$

where σ is the Stephan-Boltzmann constant, $3kT_X \sim m_X v_X^2 \sim m_X G M_X / r_X \sim \Delta E / N$ and ΔE is the energy generated in reducing r_X, the radius of the WIMP sphere containing N particles, by a factor of 2. The time required to radiate ΔE by the luminosity L_X is

$$\begin{aligned}\Delta t &= \frac{\Delta E}{L_X} \\ &= \frac{1215}{\pi^3} c^2 \hbar^3 r_X G^{-3} M_X^{-2} m_X^{-4} \\ &\sim 6.5\times 10^{13}\ \mathrm{sec}(t_A/10^7\ \mathrm{yr})^{-5/3} \\ &\quad \times (m_X/1000\ \mathrm{GeV})^{-4}(r_X/r_0)\ .\end{aligned} \tag{39}$$

Note that this mechanism requires that $r_X/l_X > 1$. The latter ratio is given by

$$\begin{aligned} r_X/l_X &= r_X n_X \sigma_{XX} \\ &\sim 0.1(\epsilon/10^{-5})(m_X/1000\ \mathrm{GeV}) \\ &\quad \times (t_A/10^7\ \mathrm{yr})^{1/3}(r_0/r_X)^2\ ,\end{aligned} \tag{40}$$

thus the WIMP sphere is opaque once $r_X \lesssim 0.3 r_0$, with r_0 being the value of r_X at the onset of the collapse.

Let us turn to neutrino cooling. A thermal neutrino population is generated when the WIMP sphere is opaque to νX scattering. The cross section for this scattering is

$$\begin{aligned}\sigma_{\nu X} &\sim \epsilon\pi^{-1}\hbar^{-4} G_F^2 (E_\nu/c^2)^2 \\ &\sim \sigma_{Xn}(3.15 kT_X/m_N c^2)^2 \\ &\sim \sigma_{Xn}(GM_X/c^2 r_X)^2 (m_X/m_N)^2\ .\end{aligned} \tag{41}$$

In the relativistic Fermi-Dirac distribution, the mean neutrino energy is 3.15 kT. Therefore,

$$\begin{aligned}\frac{r_X}{l_\nu} &\sim 5.5\times 10^{-26}(r_0/r_X)^4(\epsilon/10^{-5}) \\ &\quad \times (m_X/1000\ \mathrm{GeV})(t_A/10^7\ \mathrm{yr})^{5/3}\end{aligned} \tag{42}$$

with $l_\nu = (\sigma_{\nu X} n_X)^{-1}$ being the neutrino mean free path. The above ratio should exceed 1 for thermal neutrinos to be generated, implying

$$\begin{aligned}\frac{r_X}{r_0} &\lesssim 4.8\times 10^{-7}(\epsilon/10^{-5})^{1/4}(m_X/1000\ \mathrm{GeV})^{1/4} \\ &\quad \times (t_A/10^7\ \mathrm{yr})^{5/12}\ .\end{aligned} \tag{43}$$

The time scale for neutrino trapping is

$$\begin{aligned}\tau_\nu &\sim \frac{r_X}{c}\frac{r_X}{l_\nu} \\ &\sim 2.9\times 10^{-36}\ \mathrm{sec}(r_0/r_X)^3(\epsilon/10^{-5}) \\ &\quad \times (m_X/1000\ \mathrm{GeV})(t_A/10^7\ \mathrm{yr})^2\ .\end{aligned} \tag{44}$$

For the nominal parameters used above cooling is by Xn collisions from the beginning of the collapse and until thermal neutrinos are generated; then neutrino cooling dominates. As the contraction continues, the trapping time of the neutrinos increases while τ_{Xn} decreases and we find that in the last stages of the collapse Xn collisions dominate again. The overall cooling time is dictated by Xn scattering and is typically $\sim 10^8$ sec, very short compared to t_A, but long compared to the dynamical time scale τ_d. For larger values of m_X the second mechanism can dominate over the first one, and will take over the cooling by neutrinos in the last stages of the collapse and the overall cooling time will be even shorter than before, but still larger than τ_d. The cooling time determines the collapse time scale t_B, mentioned in Sec. I.

VII. ACCRETION OF THE NEUTRON STAR BY THE MINI BLACK HOLE

As we shall see below, the time scale for accretion of the neutron star onto the mini black hole is determined by the first phases when $M_{\mathrm{bh}} \ll M$. As the Schwarzschild radius associated with M_X, $2GM_X c^{-2} \sim 2.4\times 10^{-12}$ cm$(t_A/10^7\ \mathrm{yr})$ the problem is that of spherical accretion onto a black hole where (in the coordinates introduced in Sec. IV) the asymptotically flat spatial infinity is actually located at the star center. Therefore, the treatment of Shapiro and Teukolsky (Appendix G, Ref. 11) applies provided that the result is corrected by employing the redshift factor $(1+z_c)\sim 2$. For simplicity we adopt an equation of state

$$p = \alpha(\rho - \rho_0) \tag{45}$$

for neutron matter at densities $\gtrsim \rho_c$. From comparisons with "realistic" equations of state (e.g., Arnett and Bowers[15]) we take as representative parameters

$$\frac{\rho_c}{\rho_0} \sim 3,\quad \alpha \sim \tfrac{1}{3},\quad \rho_0 \sim 5\times 10^{14}\ \mathrm{g\,cm^{-3}}\ . \tag{46}$$

The resulting accretion rate is found to be

$$dM_{\mathrm{bh}}/dt \sim 20\pi\rho_c c(GM_{\mathrm{bh}}/c^2)^2\ . \tag{47}$$

Equation (47) is what one would expect to get for an accretion radius comparable to the Schwarzschild radius, density $\sim\rho_c$, and infall velocity $\lesssim c$. We note that, using Shapiro and Teukolsky,[11] one can derive, independently of the equation of state, a lower bound on the accretion rate which is ~ 2.5 times smaller than that of (47).

The solution of (47) gives the time scale for the consumption of the neutron star, t_C, mentioned in Sec. I:

$$t_C=(20\pi\rho_c G^2)^{-1}c^3[M_X^{-1}-M_{\rm bh}^{-1}(t_C)] \sim 1.9\times 10^5\ {\rm yr}(t_A/10^7\ {\rm yr})^{-1}\ , \quad (48)$$

where $M_{\rm bh}(t_C) \gg M_X$. Thus the star is consumed on a, relatively, very short time scale.

Can the accretion process heat up the accreting neutron matter sufficiently so that the generated neutrino luminosity will slow down the accretion? Compressional heating by itself is probably not important. For the equation of state given by Eqs. (45) and (46) we estimate that, at a radius $\sim 3GM_{\rm bh}c^{-2}$ the number density is only ~ 3 times its value at the center of the neutron star. Let us assume that some efficient dissipation mechanisms are operating in transforming the gravitational energy into thermal energy. Possible such mechanisms can be internal viscosity, or enhanced turbulent viscosity, in case the accretion flow is turbulent. Assume further that the process has a high efficiency of $\sim 10\%$. In this case the luminosity associated with the mass accretion rate of (47) is

$$L \sim 0.1 dM_{\rm bh}/dt\, c^2 \sim 4\times 10^{57}(M_{\rm bh}/M_\odot)^2 {\rm erg\,sec}^{-1}\ . \quad (49)$$

The neutrino Eddington luminosity, defined to be the luminosity for which the force exerted on a neutron by the neutrino flux equals the gravitational force, is given by[11]

$$L_{E,\nu} \sim \frac{4\pi G M_{\rm bh} c m_N}{\sigma} \sim 4.8\times 10^{57}(M_{\rm bh}/M_\odot)(E_\nu/1\ {\rm MeV})^{-2}\ {\rm erg\,sec}^{-1}\ , \quad (50)$$

where σ is the cross section for ν-n scattering $\sigma \sim \frac{1}{4}\sigma_0(E_\nu/m_e c^2)^2$ and $\sigma_0 = 1.76\times 10^{-44}$ cm^2. Even if we take the temperature, close to the horizon, as high as 100 MeV then $L < L_{E,\nu}$ for $M_{\rm bh} \lesssim 10^{-4}M_\odot$. Therefore, as long as $M_{\rm bh} \lesssim 10^{-4}M_\odot$, the accretion rate is not Eddington limited. Thus, the time scale t_C of Eq. (48) is indeed correct. We note that since the neutron star is "optically" thick for such energetic neutrinos the effective Eddington luminosity is much higher than that given in (50). What will probably occur is neutrino trapping, and diffusion outward in direction and downward in energy.

After the mini black hole reaches a mass of $\sim 10^{-3}M_\odot$, it will accrete the neutron star on a time scale of ~ 10 sec. One may expect a neutrinosphere with an initial temperature ~ 10 MeV, and an average neutrino energy which is decreasing during these last few seconds due to cooling and increasing redshift. When the neutrinosphere approaches the black-hole horizon, the neutrino signal will terminate abruptly. The surface of the infalling neutron star will also radiate MeV photons with average energy that, also, decreases with time. The characteristics of this γ-ray emission (peak energy, duration, and spectral softening with time) are similar to those of γ-ray bursters.[19] However, there are indications that the latter are probably recurrent,[20] unlike in the present case.

VIII. DISCUSSION

We examined the characteristics of those WIMP's that can destroy neutron stars on a time scale shorter than observed ages of neutron stars. These WIMP's are therefore ruled out as candidates for the galactic missing mass. The ruled out WIMP's were found to have rather special characteristics. They are either bosons of mass 10^5 GeV $\gtrsim m_X \gtrsim 200$ GeV that however do not have coherent vectorial repulsive interactions or, either bosons or fermions in this mass range, that have a strong enough coupling to an attractive coherent scalar that overcomes the repulsive vector interaction as well as the Fermi degeneracy pressure. The upper limit on m_X follows from considerations of annihilation of $X\bar{X}$ in the early Universe while the lower bound is discussed below. For bosons, Bose condensation pushes the above lower limit further down to ~ 1 GeV. As discussed in Sec. IV, for bosonic WIMP's of mass 1 GeV $\lesssim m_X \lesssim 1.7\times 10^4$ GeV, Bose condensation determines the time scale for the onset of self-gravity. In the following we consider the more conservative lower limits on m_X resulting from (25), which does not rely on the occurrence of Bose condensation.

The derived constraints are obtained from observations of neutron stars in various astrophysical systems. We have found that the overall time scale is determined by $t_A(\propto T^{3/2})$, with T being the interior temperature of the neutron star. Thus, we should look for neutron stars that are as cold and as old as possible. Cooling curves of neutron stars[16,17] imply that, in the absence of heating, the interior temperature of an isolated neutron star gets down to 10^5 K in $\sim 10^7$ yr and to 10^4 K in $\sim 3\times 10^7$ yr. However, even isolated neutron stars probably cannot get below $\sim 10^5$ K due to heating by accretion of interstellar matter.[17] Other possible sources of heating are due to glitches and to viscous dissipation in the neutron star matter.[17] In accreting neutron stars that are members of a binary system the surface temperature and hence the interior temperature could get as high as few times 10^7 K.

Solitary radio pulsars can be observed only if younger than $\sim 10^7$ yr, because of the decay of the magnetic field.[21] Assuming an interior temperature of 10^5 K we find that WIMP's, with the above special characteristics, of masses exceeding $\sim 2\times 10^4$ GeV will destroy neutron stars on a time scale shorter than the observed ages. For a higher WIMP density, corresponding to locations closer than the Sun to the galactic center, the lower limit on m_X can be somewhat reduced.

Consider next neutron stars which are members of galactic bulge or globular clusters weak x-ray sources (luminosity $\lesssim 10^{34}$ ergs sec^{-1}). The surface temperature of these stars is $\sim 10^6$ K, and taking it to be also the interior temperature yields $t_A \lesssim 10^9$ yr, for $m_X \gtrsim 10^4$ GeV. The estimated ages are $\gtrsim 10^9$ yr (Refs. 22 and 23). Thus, WIMP's with the special characteristics mentioned above, of $m_X \gtrsim 10^4$ GeV are ruled out.

Another interesting class of objects is that of low-mass compact x-ray binaries in which it is believed that the neutron star evolved out of a white dwarf that exceeded its Chandrasekhar mass limit due to accretion from its companion.[22,23] The analogue of (19) for a white dwarf gives an accretion rate ~ 230 times larger than that for a neutron star. As we mention below, the WIMP's captured in a white dwarf will not become self-gravitating. Thus, the white dwarf phase will serve only as a collecting phase for the subsequent neutron-star phase. While the interior temperature in these neutron stars is high compared to that of an isolated pulsar (few times 10^7 K), the increased value of $M_X(t_A)$ combined with the large age of these systems ($>10^9$ yr) (Refs. 22 and 23) can rule out masses $m_X \gtrsim 2\times 10^3$ GeV.

Millisecond pulsars have estimated ages $\gtrsim 10^9$ yr (Refs. 24–27) and probably temperatures as low as a few times 10^5 K. The resulting lower bound on m_X is $\sim 800-1500$ GeV. It was suggested recently[28] that the neutron stars in the millisecond pulsars formed from a white dwarf. In this case, even for interior temperature of 10^6 K, we get $m_X \gtrsim 200$ GeV.

Can the above outlined scenario apply also to white dwarfs? Repeating the procedure of Sec. III, for the case of a white dwarf, we find that for a typical white dwarf of mass $1.1 M_\odot$, and radius 5×10^8 cm (Ref. 29), Eq. (18) now gives $\epsilon \gtrsim 0.08$ and $m_X \lesssim 2.2\times 10^4 \epsilon$ GeV. For WIMP's satisfying these constraints, Eq. (19) yields an accreted mass flux of $\sim 3.6\times 10^{18}$ g$(t/10^7$ yr). For a central density of $\sim 5\times 10^7$ g cm^{-3} (Ref. 29) and central temperature of $\sim 10^7$ K (Ref. 11), one gets from Eq. (25) $t_A \sim 1.4\times 10^{13}$ yr. Therefore, the WIMP's will not become self-gravitating and a mini black hole will not form. However, as mentioned above, in cases that the white dwarf evolves into a neutron star the latter inherits the accumulated WIMP mass of the preceding white-dwarf phase. We note once more that these WIMP's must have relatively low masses and large ϵ. A Bose condensation in a white dwarf cannot happen as it requires, for the above parameters, that $T \lesssim 340$ K$(t/10^7$ yr$)^{1/3}(m_X/1000$ GeV$)^{-1/3}$, which is orders of magnitude lower than the interior temperature of a white dwarf.

If the WIMP's do not possess the characteristics mentioned above, they would still be trapped by neutron stars and drift toward its center as they cool down. They would even become self-gravitating, but the collapse to a mini black hole would be prevented either by Fermi degeneracy pressure or by a repulsive vectorial coherent interaction. These mass concentrations will represent only a slight (typically $\sim 10^{-12}$) perturbation to the gravitational potential at the center of the neutron star and will have only negligible effect on its structure. Recently, Eichler[30] considered the possibility that exotic matter hidden in main-sequence stars can get excavated by various astrophysical processes and reveal itself as compact ultralight satellites or binary companions. In the case of neutron stars, a disruption of the neutron star is a necessary (but probably insufficient) condition for the excavation of the WIMP sphere. Disruption of neutron stars could occur in binaries where both members are neutron stars (e.g., the binary pulsar PSR1913+16). In these systems, the orbital separation decays due to emission of gravitational radiation leading eventually to tidal disruption of the less massive neutron star.[31]

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