White Dwarfs as Dark Matter Detectors

Ryan Janish, Vijay Narayan, and Paul Riggins
Berkeley Center for Theoretical Physics, Department of Physics,
University of California, Berkeley, CA 94720, USA

White dwarfs can serve as detectors for ultra-heavy dark matter states which interact to trigger type Ia supernovae. This was originally proposed in [1] and used to place bounds on primordial black holes. In this paper we extend the capability of white dwarf detectors to candidates with non-gravitational couplings, focusing on dark matter transits and collisions within the white dwarf. In particular, we provide a detailed analysis of the explosiveness for any heating mechanism in the white dwarf which releases high-energy standard model particles. We apply this mechanism to constrain Q-ball dark matter and model of dark nuclei in regions of parameter space fundamentally inaccessible to terrestrial-based experiments.

PACS numbers:

I. INTRODUCTION

The detection of ultra-heavy dark matter (DM) is an open problem which will ultimately require a confluence of astrophysical probes. For instance, DM masses above $\sim 10^{22}$ GeV will register fewer than 1 event per year in a typical terrestrial detector of size $\sim (100 \text{ m})^2$. Super-K has SA $\sim 7.5 \times 10^7 \text{cm}^2$. In addition, the lack of conclusive signatures on a variety of experimental fronts has led many to seriously consider DM candidates far above the weak scale and their potential signatures citations. One possibility proposed by [1] is that ultra-heavy DM can trigger supernovae in sub-Chandrasekhar white dwarf stars (WD) by inducing runaway fusion. In this regard, white dwarfs can serve as detectors for ultra-heavy DM states.

White dwarfs are particularity suited to this task, as they are drastically more susceptible to runaway fusion than are main-sequence stars. Runaway fusion requires that two criteria be met: a region within the star must be hot enough to support exothermic fusion reactions, and the rate at which energy is released by these reactions must dominate any cooling mechanisms that drain energy from the fusing region. The stellar medium of a WD has fewer cooling mechanisms available than does a non-degenerate star - WD cooling relies on thermal diffusion whereas main-sequence stars can also cool via thermal expansion. This is suppressed in a WD because its pressure is determined by electron degeneracy and is thus independent of temperature. The remaining primary cooling mechanism, diffusion, becomes less important over longer length scales and can be overcome by sufficiently heating a large enough region of the WD.

The necessary trigger for runaway fusion was initially computed in [2] and subsequently implemented in [1] to constrain primordial black holes, which can ignite the star via gravitational dynamical friction. In addition, the authors of [1] identify several other heating mechanisms involving DM which may be constrained in a similar manner. Some of these have been explored in follow-up works...citations, citations. In this work, we extend their analysis to DM candidates with generic non-gravitational couplings, focusing on DM transits through the WD and DM-DM collisions within the WD which release energy in the form of standard model (SM) particles. Concrete examples of DM candidates with interactions of this type include baryonic Q-balls found in supersymmetric extensions of the SM and dark nuclei with high-order couplings to the SM. We are able to constrain large regions of parameter space for these models, and do so in regions fundamentally inaccessible to terrestrial experiments. More generally, we provide a detailed analysis of the explosive power and resulting constraints on any DM with interactions of this sort. Other ignition sources exist (i.e., detector backgrounds). How does this limit our constraints? Discuss the specific tests we apply (WD lifetime, SN rate) to derive our constraints.

We should emphasize somewhere that our constrains are by nature complimentary to terrestrial ones - it is more massive DM that is likely to cause boom, and also more massive DM that is hard to see on Earth due to low flux. We also have a "flux is too low" limit, but at much higher masses. Might be worth giving this mass explicitly here and comparing to the analogous terrestrial detector limit.

Add the traditional table-of-contents paragraph

Possible long term addition: improve constraints by using non-type IA sn rate instead of total rate. Should talk with SN people first.

II. REVIEW: WHITE DWARF RUNAWAY FUSION

In a WD, the *fusion temperature* T_B is a constant set by the energy required for ions to overcome their mutual Coulomb barriers $T_B \sim$ MeV. WD cooling is set by the thermal diffusivity of photons and degenerate electrons. The two species dominate the cooling at different stellar densities, as determined precisely in [2]. Note that for a heated region of size R, this cooling rate scales as R^{-2} while the fusion rate scales as R^{-3}. Thus there is a critical *trigger size* λ_T below which diffusive cooling dominates the thermal evolution of a temperature peak, and above which the liberated fusion energy dominates [2]. Therefore, a region with temperature greater than T_B and size greater than λ_T in a given WD will launch a runaway fusion chain-reaction and result in a type Ia supernovae. The numerical value of λ_T is highly sensitive to the WD density and has been analytically scaled for varying WD masses in [1]. These results are reproduced in Figure 1. As in [1], we restrict our attention to carbon-oxygen WDs in the upper mass range $0.7 - 1.4\ M_\odot$ which correspond to a number density of nuclei $n \sim 10^{29} - 10^{32}\ \text{cm}^{-3}$. Over this range, the trigger size is approximately $\lambda_T \sim 10^{-5} - 10^{-2}$ cm.

A general energy deposition event in the WD can be roughly characterized by two parameters: a temperature T and heating length L. This energy deposit will eventually take the form of a local peak in the WD temperature profile. We define T to be the characteristic temperature and L the characteristic length scale of this local peak as it initially appears. L is determined by the efficiency with which a given energy deposition mechanism interacts with the stellar medium. For instance, suppose that kinetic energy were transferred directly to neighboring ions via elastic scatters. These ions would thermalize over their collisional time scale resulting in a heating length L of order the ion mean free path. In fact, this is the minimal L possible for any process which couples to WD constituents (ions or electrons). In the other extreme, suppose that a process produces a large number of electrons with energy just above the Fermi energy. These electrons have Pauli-suppressed interactions with the medium and will travel a long distance before their energy is scattered and thermalized, resulting in a large L - possibly of order the stellar radius.

We find that a given heating event will destroy the WD as long as the following is satisfied:

$$T \gtrsim T_B, \qquad L \gtrsim \lambda_T. \tag{1}$$

Note that this condition determines if a region of size L and temperature T *immediately* initiates runaway fusion. Instead, it is more useful to determine whether a given energy deposit will *eventually* initiate runaway fusion. In particular, a heating event that results in $T \gg T_B$ and $L < \lambda_T$ is perhaps initially dominated by diffusive cooling but can potentially evolve into an "explosive" temperature profile capable of igniting the star. When diffusion dominates, the thermal evolution will approximately be given by the dilution of the fixed total energy of the initial region over the final volume. Taking this into account, we can express the explosion condition (1) in terms of an excess energy within the temperature peak (rather than the temperature itself):

$$E \gtrsim n T_B \text{Max}\left[L, \lambda_T\right]^3. \tag{2}$$

The moral of (2) is that for a heating event to be capable of destroying a certain WD, both the energy density and the energy deposited must be sufficiently large. In fact, there is an absolute minimum explosion energy given by

$$E_{\text{boom}} \sim n T_B \lambda_T^3 \sim 10^{14} - 10^{22}\ \text{GeV}, \tag{3}$$

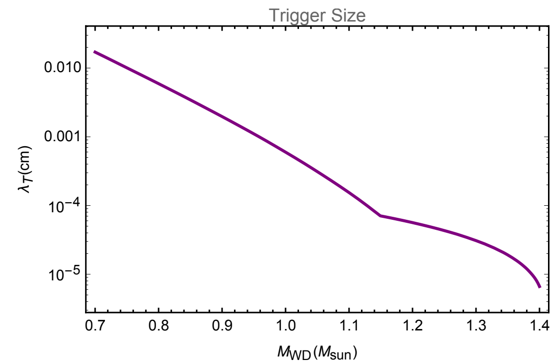


FIG. 1: Approximate trigger sizes for carbon-oxygen white dwarfs, based on numerical results in [2] and analytically scaled in [1]. extend to 0.4 solar mass

where λ_T varies over a range of WD densities as shown in Figure 1. However, this energy is sufficient only if it thermalizes within the trigger size. If an energy is deposited on a length scale larger than λ_T, the explosion energy required is given by (2) and may be parametrically larger. Since we are concerned with processes that have a fixed deliverable energy (the energy of the incoming DM), the most explosive processes will be those that result in the most localized heating. As a result, understanding the heating length L for various processes is critical to assessing their explosive potential. This is done in Section IV.

III. DARK MATTER EXPLOSIVENESS

In this section, we parameterize the explosive power of a generic (ultra-heavy) DM interaction in the WD which releases n_i SM particles of species i each with kinetic energy ϵ.

A. DM Transit

Consider a DM transit through the WD, interacting with stellar constituents (ions or electrons) in a general manner as shown in Figure III A. The cross section for this interaction is denoted as $\sigma_{i,\epsilon}$. Of course, this released energy must be transferred to the stellar medium in order for the WD to be heated. For a given particle type, each value of ϵ is characterized by a distance $R(\epsilon)$ from the point of release over which it is deposited. In particular, we define $R(\epsilon)$ as the distance over which a particle i and any secondaries transfer $\mathcal{O}(1)$ of the initial energy ϵ to electrons or ions in the WD. This will ultimately be related to the heating length L.

Consider the two possibilities relevant for ignition. If $R_\epsilon > \lambda_T$, the DM must deposit a minimum energy $E_{\text{boom}} \sim R_\epsilon^3 n T_B$ in order to heat up the entire region of size R_ϵ to the critical temperature T_B, where n is the

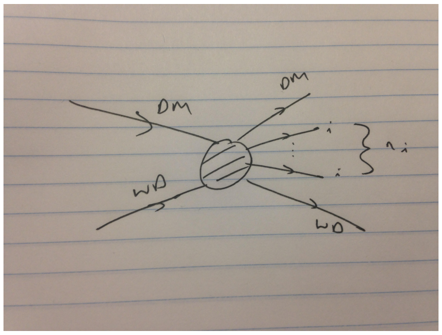


FIG. 2: General interaction between ultra-heavy DM and WD constituents, producing n_i additional particles of species i.

number density of nuclei in the WD. On the other hand, if $R_\epsilon < \lambda_T$ then the minimum energy required is independent of R_ϵ and given by $E_{\rm boom} \sim \lambda_T^3 n T_B$. Setting $T_B \sim$ MeV, $\lambda_T \sim 10^{-5}$ cm, and $n \sim 10^{32}$ cm^{-3}, we find that an energy $E_{\rm boom} \sim 10^{14}$ GeV transferred to the WD within a localized region smaller than λ_T^3 will eventually trigger runaway fusion. This "explosion" energy must be less than the total energy released during the DM transit. Doing so, we find a lower bound on the interaction cross section sufficient to trigger runaway fusion:

$$n_i \sigma_{\epsilon,i} \gtrsim \begin{cases} \lambda_T^2 \left(\dfrac{T_B}{\epsilon}\right) & R_\epsilon < \lambda_T \\ \lambda_T^2 \left(\frac{R_\epsilon}{\lambda_T}\right)^2 \left(\frac{T_B}{\epsilon}\right) & R_\epsilon > \lambda_T \end{cases}. \tag{4}$$

Note that in deriving this explosiveness bound, we have assumed the DM transit time is less than the corresponding diffusion time. After a time Δt the DM has traversed $v_{\rm esc}\Delta t$, where the velocity is set by the escape velocity of a WD. Therefore, this amounts to the following condition:

$$\begin{aligned} &\tau_d(\lambda_T, T_B) \gtrsim \tfrac{\lambda_T}{v_{\rm esc}}, \qquad R_\epsilon < \lambda_T \\ &\tau_d(R_\epsilon, T_B) \gtrsim \tfrac{R_\epsilon}{v_{\rm esc}}, \qquad R_\epsilon > \lambda_T \end{aligned} \tag{5}$$

where τ_d is the characteristic time for a region of given size and temperature to diffuse $\mathcal{O}(1)$ of its heat. An estimate from the heat equation gives $\tau_d(\lambda_T, T_B) \sim \frac{\lambda_T^2}{\alpha}$, where α is the (temperature-dependent) diffusivity. Show condition is true for all densities. We also assume that the time to transfer energy ϵ out to its characteristic range R_ϵ is less than the diffusion time scale $\tau_d(\lambda_T, T_B)$.

B. DM-DM Collisions

IV. DETERMINATION OF HEATING LENGTH L

In this section, we enumerate the different ways in which standard model particles can lose energy to the WD (unless otherwise noted, we will assume a carbon-oxygen WD.) The interior of a WD is a complex environment. Famously, the star is supported against collapse by electron degeneracy pressure with a characteristic Fermi energy $E_F \sim n_e^{1/3} \sim \mathcal{O}$(MeV) at all densities?. In addition, the nuclei are at an ambient temperature $T \sim$ keV and form a strongly-coupled plasma $\frac{Ze^2}{n^{1/3}T} \gg 1$. In what follows, we calculate the WD medium sensitivity to deposited energy and the deposit ranges R_ϵ for the different possible DM interactions. For the purpose of depositing sufficient energy to trigger supernovae, we focus on high-energy particles $\epsilon \gg$ MeV which interact via the strong and electromagnetic forces (i.e. electrons, muons, photons, pions, and neutral hadrons). In this way, the WD may be thought of as a "particle detector" with electromagnetic and hadronic "calorimeter" components.

A. Electromagnetic Interactions

For charged particles, Coulomb scattering is a useful mechanism for energy transfer. Generically, an incident (spin-0) particle of mass m_i, charge e, and velocity β scattering off a target M_t of charge Ze is described by the "Rutherford" differential cross section Originally derived by Bhabha?

$$\frac{d\sigma}{dE'} = \frac{2\pi\alpha^2 Z^2}{M_t \beta^2} \frac{1}{E'^2} \left(1 - \frac{\beta^2 E'}{E_{\rm kin}}\right), \tag{6}$$

where we have assumed a sufficiently fast incident particle so that interactions are governed by single collisions with energy transfer E' [3]. $E_{\rm kin}$ denotes the maximum energy transfer possible satisfying kinematic constraints target at rest and zero relative angle between incoming incident and outgoing target momenta:

$$E_{\rm kin} = \frac{2M_t \beta^2 \gamma^2}{1 + 2\gamma(M_t/m_i) + (M_t/m_i)^2}, \tag{7}$$

where $\gamma = (1-\beta)^{-1/2}$.

For sufficiently heavy incident particles, the differential cross section depends only on the velocity of the incident particle. Note that higher-spin particles receive additional corrections to the cross section, but for small energy transfers these corrections are negligible. It is straightforward to understand the parametric dependences of (6): there is increased likelihood to scatter for slowly moving incident particles undergoing "soft-scatters" against lighter targets. Therefore, one would expect that soft scattering dominates the energy loss and that collisions with nuclei of mass M are suppressed by a factor $\mathcal{O}\left(\frac{Zm_e}{M}\right)$ as compared to collisions with electrons. This is certainly true for incident charged particles in non-degenerate matter. However, both of these naive expectations turn out to be false when considering scattering off a degenerate species.

To understand the effect of degeneracy, we first consider the energy loss from scattering a high-energy charged particles off non-degenerate targets in the WD, for instance, the carbon nuclei. In this case, the stopping power due to collisions with a number density n is given by:

$$\frac{dE}{dx} = -\int dE' \left(\frac{d\sigma}{dE'}\right) nE' \tag{8}$$

$$\sim -\frac{2\pi n Z^2\alpha^2}{M_t\beta^2}\log\left(\frac{E_{\max}}{E_{\min}}\right). \tag{9}$$

This integration must be performed over all E' within the regime of validity for (6), fixing the lower and upper bounds of the "Coulomb Logarithm". Quantum mechanical uncertainty sets a limit to the accuracy that can be achieved in "aiming" an incident particle at a target. In terms of impact parameter b for the collision, this translates to a bound $b \gtrsim (\min\{m_i, M_t\}\beta\gamma)^{-1}$ corresponding to the larger de Broglie wavelength. In terms of energy transfer, this becomes

$$E' \lesssim E_q = \frac{2\,\min\{m_i, M_t\}^2 Z^2\alpha^2\gamma^2}{M_t}. \tag{10}$$

In addition, the expression for the differential cross section (6) no longer holds when E' becomes larger than the mass of the target [4]. For our calculations, we take the maximum energy that an incident particle is able to transfer to be

$$E_{\max} = \min\{E_q, E_{\rm kin}, M_t\}. \tag{11}$$

On the other hand, the maximum impact parameter is determined by charge screening. In a WD, this is simply the screening due to a degenerate electron gas similar to a solid and is given by the Thomas-Fermi length [5]

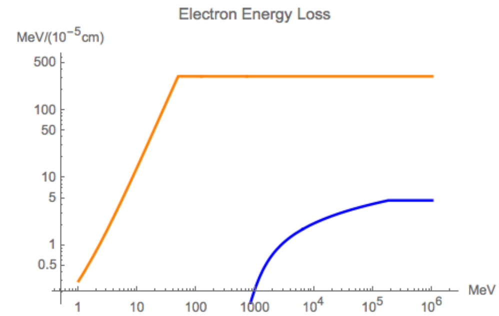
$$l_{\rm sc} = \left(\frac{6\pi Z e^2 n_e}{E_F}\right)^{-1/2} \sim \frac{1}{m_e}. \tag{12}$$

This is the equivalent of the Debye screening length for a degenerate gas at Fermi energy E_F. This corresponds to a lower bound

$$E_{\rm sc} = \frac{2m_e^2 Z^2\alpha^2}{M_t\beta^2}. \tag{13}$$

Note that when the incident particle reaches velocity $\beta\gamma \approx \frac{m_e}{\min\{m_i, M_t\}}$, the minimum possible energy transfer due to Thomas-Fermi screening will in fact exceed the maximum. At this point, the form of equation (8) becomes modified to avoid unphysical non-negative values of dE/dx and there is negligible stopping power due to collisions. However, the lattice structure of ions in the WD introduces further complications [5]. The typical lattice binding energy between ions is given by the electric potential

$$E_B \sim \frac{Z^2 e^2}{n^{-1/3}}. \tag{14}$$



For energy transfers greater than E_B, any lattice effects can be safely ignored. On the other hand, momentum transfers from collisions below this threshold may lead to suppressed energy loss due to collective effects. For simplicity we set the lower bound on nuclei scattering (not applicable for electron targets) to be

$$E_{\min} = \max\{E_B, E_{\rm sc}\} \tag{15}$$

so that we only account for energy transfers greater than the ionic lattice binding energy.

When considering collisions with the degenerate electrons, an incident particle transferring energy E' can only scatter those electrons within E' of the Fermi surface. We define a modified density of electrons $n_e(E')$ as:

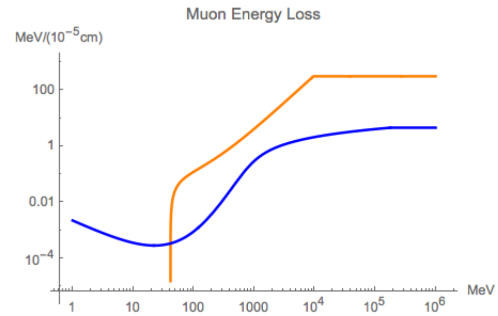
$$n_e(E') = \begin{cases} \displaystyle\int_{E_F - E'}^{E_F} dE\ g(E) & E' \leq E_F \\ n_e & E_F \leq E' \end{cases}, \tag{16}$$

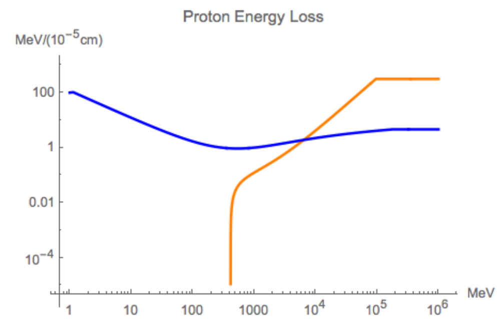
where $g(E)$ is the density of states per unit volume for a three-dimensional free electron gas. This can also be expressed as a suppression of the differential cross section of order $\mathcal{O}(E'/E_F)$ whenever energy less than E_F is transferred. Therefore, unlike in the non-degenerate case, the energy loss due to soft-scatters are in fact subdominant to the contributions from rare hard-scatters. The stopping power is also highly sensitive to $E_{\max}$ and $E_{\min}$, governed by a power-law dependence rather than the logarithmic sensitivity in the case of a non-degenerate target.

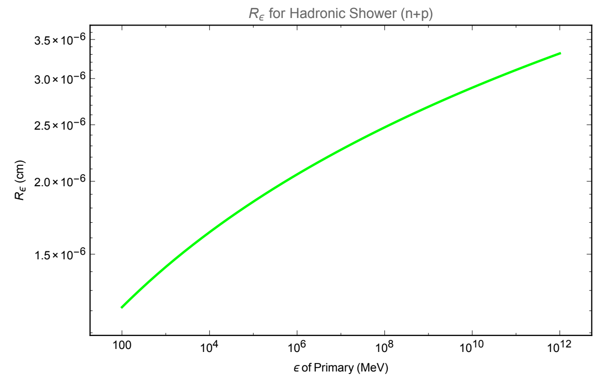
We have calculated the stopping power of high-energy particles solely due to Coulomb collisions, differentiating between nuclei and degenerate electron targets. This is shown in Figure Figure for an electron density $n_e \sim 10^{33}\ \text{cm}^{-3}$. It is found that for heavy incident particles, the stopping power is dominated by collisions with nuclei at low-energies although it is dominated by collisions with degenerate electrons at high-energies. In addition, scattering off degenerate electrons becomes completely screened at a comparatively higher energies.

B. Strong Interactions

Strong interactions are a significant channel for energy transfer in the WD. For sufficiently energetic par-







ticles greater than the nuclear binding energy $E_{nuc} \sim \mathcal{O}(10 \text{ MeV})$, nuclear interactions (absorption) result in an $\mathcal{O}(1)$ number of energetic secondary hadrons (protons, neutrons, pions, etc.) emitted roughly in the direction of the primary particle. These secondary particles approximately split the initial total energy of the absorbed primary particle and can collide with other nuclei in the WD. In addition, the nucleus will generally be left in an excited state and relax through the emission of low-energy $\mathcal{O}(10 \text{ MeV})$ nucleons ("nuclear evaporation") and photons [4]. The cross section for any such process is approximately set by the nuclear length scale $\sim$ fm. A hadronic shower is a result of all such reactions caused by primary and secondary particles.

To estimate the typical shower length, we assume a primary particle of energy ϵ interacting with nuclei to produce N secondary particles each of energy ϵ/N. The shower will end once final-state particles reach energies of order E_{nuc}. We ignore the effects of the evaporative stage in each reaction as this will only introduce minor corrections at sufficiently large ϵ. We find a hadronic shower length $L = l_{\rm nuc}\left(\frac{\log(\epsilon/E_{nuc})}{\log N}\right)$, where $l_{\rm nuc}$ is the mean free path for nuclear collisions. For a number density $n \sim 10^{32}$ cm^{-3} and hadronic cross section ~ 0.1 barn match nuclear data for carbon nonelastic scattering this results in a shower length $L \approx 10\ l_{\rm nuc} \approx 10^{-6}$ cm. Note that this only has a mild, logarithmic sensitivity to the initial energy $\epsilon \gg$ 10 MeV and number of secondaries $N \sim \mathcal{O}(1)$.

The end products of a hadronic shower will generally be hadrons of kinetic energy $1-10$ MeV which are incapable of inducing further nuclear disintegration. Charged hadrons have appreciable electromagnetic interactions with nuclei, while collisions with degenerate electrons are completely screened at these energies. Integrating (8), we find that in a density of $n \sim 10^{32}$ cm^{-3} protons traverse $\approx 10^{-6}$ cm slowing down from 5 MeV to 1 MeV. The corresponding distance is significantly larger for charged pions $\sim 10^{-2}$ cm On the other hand, electromagnetic couplings are highly suppressed for final-state neutral hadrons. The dominant source of energy loss for neutrons in this range is elastic scattering off nuclei neutron capture cross section in C+O is negligible. Each scatter transfers a fraction $\left(1-\left(\frac{m}{m+M}\right)^2\right) \approx 0.1$ of the neutron energy to nuclei (averaging over scattering angles), where m and M are the masses of the neutron and nuclei, respectively. At kinetic energies below $\sim$ MeV, the typical energy transfer in an elastic collision is of order the binding energy in the Coulomb lattice E_B. We find that $\mathcal{O}(10)$ elastic scatters are needed to slow neutrons from 5 MeV to 1 MeV. For a number density $n \sim 10^{32}$ cm^{-3} and elastic cross section $\sim$ barn set to match nuclear data for carbon nonelastic scattering, we find that neutrons traverse a total distance (in the form of a random-walk) $\sim 10^{-8}$ cm during this energy deposition.

V. EXAMPLE: Q-BALL DARK MATTER

In various supersymmetric extensions of the standard model (SM), non-topological solitons called Q-balls can be produced in the early universe [6, 7]. If these Q-balls were stable, they would comprise a component of the dark matter today. Q-balls can be classified into two groups: supersymmetric electrically charged solitons (SECS) and supersymmetric electrically neutral solitons (SENS). When a neutral baryonic Q-ball interacts with a nucleon, it absorbs its baryonic charge as a minimum-energy configuration and induces the dissociation of the nucleon into free quarks. In this process (known as the "KKST" process), $\sim$ GeV of energy is released through the emission of 2-3 pions [8]. The KKST process provides a useful way to detect such Q-balls. The cross section for interaction is approximately the geometric cross section

$$\sigma_Q \simeq \pi R_Q^2. \tag{17}$$

In gauge-mediated models with flat scalar potentials, the Q-ball mass and radius are given by

$$M_Q \sim m_F Q^{3/4}, \quad R_Q \sim m_F^{-1} Q^{1/4}, \tag{18}$$

where m_F is related to the scale of supersymmetry breaking (messenger scale). The condition $M_Q/Q < m_p$ ensures that the Q-ball is stable against decay to nucleons [8].

Note that a sufficiently massive Q-ball will become a black hole if the Q-ball radius is less than the Schwarzschild radius $R_Q \lesssim R_s \sim GM_Q$. In the model described above, this translates into the condition

$$m_F \left(\frac{M_{\rm pl}}{m_F}\right)^3 \lesssim m_Q, \quad \left(\frac{M_{\rm pl}}{m_F}\right)^4 \lesssim Q. \tag{19}$$

For Q-ball masses of this order, gravitational interactions become relevant while the KKST interaction ceases to exist.

A. Q-ball Explosiveness

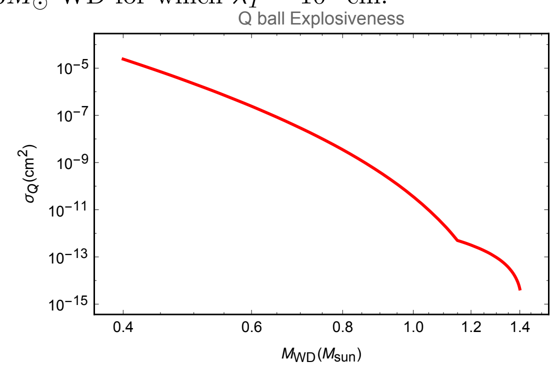
We assume that for each Q-ball collision, there is equal probability to produce π^0, π^+ and π^- under the constraint of charge conservation. In the notation of Section overview, $n_\pi \sim \mathcal{O}(10)$ pions are released with $\epsilon_\pi \sim$ GeV. The mean distance travelled by a relativistic particle before decaying is $d = \gamma v \tau$. $\gamma \approx 5$, so for neutral pions $d_{\pi^0} \sim 10^{-5}$ cm while for charged pions, $d_{\pi^\pm} \sim 10$ m.

Numerous experiments have studied the effects of $50 - 500$ MeV pions incident upon complex nuclei targets such as carbon. It is found that there is roughly equal cross section of order ~ 0.1 barn for a (neutral or charged) pion to either scatter elastically, scatter inelastically, or become absorbed with no final state pion [9]. Of these possibilities, pion absorption is the most relevant for energy loss as this will induce a hadronic shower. During this process, $\sim$ 2-3 hadrons (i.e. protons, neutrons, alpha particles) are emitted. Point of this paragraph: argue that KKST range lies within trigger size

Therefore, Q-balls with sufficiently large cross section $\sigma_Q \gtrsim 10^{-4}\lambda_T^2$ will result in explosions. This is plotted in Figure Figure. For the model described in Section previous Section, the Q-ball cross section is related to its mass $m_Q = \sigma_Q^{3/2} m_F^4$ and $Q = \sigma_Q^2 m_F^4$. Therefore, find that

$$m_Q \gtrsim 10^8 \text{ g} \left(\frac{m_F}{\rm TeV}\right)^4, \quad Q \gtrsim 10^{38} \left(\frac{m_F}{\rm TeV}\right)^4 \tag{20}$$

is capable of triggering runaway fusion in a heavy $\sim 1.25 M_\odot$ WD for which $\lambda_T \sim 10^5$ cm.



B. Q-ball Constraints

The transit rate for an ultra-heavy DM state through a WD is given by

$$\Gamma_{\rm DM} = n_Q \sigma_g v \sim \frac{\rho_{\rm DM}}{m_{\rm DM}} \pi R_s^2 \left(\frac{v_{\rm esc}^2}{v}\right), \tag{21}$$

where $v \sim 10^{-3}$ is the virial velocity of DM and $\rho_{\rm DM}$ is the energy density of DM in the region of interest. σ_g denotes the capture cross section including the gravitational Sommerfeld enhancement.

The existence of heavy white dwarfs can rule out Q-ball DM. Considering a $1.25 M_\odot$ WD in the local dark matter halo $\rho_{\rm DM} \sim 0.3$ GeV/cm^3, we find that a Q-ball of mass $m_Q \lesssim 10^{20}$ g will transit with a rate $\Gamma_Q \gtrsim$ Gyr^{-1}. If we instead consider (recently discovered) heavy white dwarfs in the galactic center $\rho_{\rm DM} \sim 10^3$ GeV/cm^3, Q-balls of mass $m_Q \lesssim 10^{24}$ g will transit within a Gyr.

VI. DISCUSSION

Appendix A: Detail Study of Heating Lengths

Appendix B: Acknowledgements

We are grateful to S. Rajendran for suggestions and collaboration during the early stages of this work, as well as his "benign negligence". We would also like to thank D. Grabowska, K. Harigaya, R. McGehee, and J. Wurtele for stimulating discussions.

[1] P. W. Graham, S. Rajendran, and J. Varela, Phys. Rev. **D92**, 063007 (2015), 1505.04444.

[2] F. X. Timmes and S. E. Woosley, Astro. Phys. Journal **396** (1992).

[3] K. A. Olive et al. (Particle Data Group), Chin. Phys. **C38**, 090001 (2014).

[4] Bruno Rossi, *High Energy Particles* (Prentice-Hall, Inc., Englewood Cliffs, NJ, 1952).

[5] Stuart L. Shapiro and Saul A. Teukolsky, *Black Holes, White Dwarfs, and Neutron Stars* (Wiley, 1983).

[6] S. R. Coleman, Nucl. Phys. **B262**, 263 (1985), [Erratum: Nucl. Phys.B269,744(1986)].
[7] A. Kusenko and M. E. Shaposhnikov, Phys. Lett. **B418**, 46 (1998), hep-ph/9709492.
[8] M. Dine and A. Kusenko, Rev. Mod. Phys. **76**, 1 (2003), hep-ph/0303065.
[9] T. Lee and R. Redwine, Annu. Rev. Nucl. Part. Sci. **52** (2002), 1007.0039.