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Charged Dark Matter

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We propose that dark matter is made of CHAMPs, charged massive particles that survive annihilation in the early Universe. We constrain champs to be matter-antimatter symmetric, to provide critical mass density, and to constitute the non-baryonic halo of galaxies such as ours. We show that the window of allowed champ mass extends from 20 to 1000 TeV. Champs of charge +1 should now appear as super-heavy isotopes of Hydrogen. While some of their antiparticles could bind to ^4He nuclei somewhat after primordial nucleosynthesis (also to appear today in the disguise of super-heavy Hydrogen), we argue that negative champs overwhelmingly bind to protons to pose as super-heavy stable neutrons. *En route* to detectors, these bound champs suffer nuclear reactions that change their ultimate appearance. By the time they come to rest, negative champs have recombined with larger nuclei to form various super-heavy isotopes. We discuss how and where to search for relic champs.

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1. Introduction

In lustra past, particle theorists explained particles that were known to exist (like pion-nucleon resonances and strange particles) and predicted others that had to exist (like charmed particles, W 's and Z 's). Overwhelmed by the successes of the standard model, they now find themselves all too often enumerating the properties of particles that have no reason *not* to exist. This is especially evident in countless considerations of candidates for the dark matter of galactic halos and the missing matter of the Universe, whose identification is surely a central problem at the frontier of cosmology and particle physics. In the present endeavor, we introduce a seemingly outrageous dark matter candidate that fulfils its task admirably nonetheless: a stable and very massive electrically-charged particle.

Angular-momentum conservation and electromagnetic gauge invariance coerce photons, electrons and the lightest neutrinos to be absolutely stable. Other particles, including protons, many atomic nuclei, and (perhaps) the heavier neutrinos, have lifetimes longer than the age of the Universe. In the context of the standard model, their stability stems from shallow cause, since it is simply not possible to construct renormalizable interactions to make these particles decay rapidly, and it is all too easy to impose *ad hoc* symmetries to make them absolutely stable. Most matter we know of is stable for reasons we cannot yet comprehend. Under these circumstances, what fool would argue that photons, electrons, neutrinos, protons and bound neutrons must exhaust the list of particles that have survived since the Big Bang? The problem in finding suitable dark-matter candidates is the embarrassment of riches.

Aside from nuggets of strange quark matter and objects of astronomical proportion, the usual list of dark-matter candidates includes only neutral particles and sparticles. We examine the novel possibility that dark matter consists of elementary particles of nonvanishing integral electric charge, hypothetical particles we call CHAMPs and denote by $C^\pm$. To have avoided detection, champs must be very massive. Under earthly conditions, charged champs would be — or would be bound to — atomic nuclei, depending on whether their charge is positive or negative. With one exception, champs disguise themselves as preposterously heavy isotopes of known chemical elements. Negative champs bound to protons, however, correspond to the missing zeroth entry in the Periodic Table. They play an especially important rôle in our cosmology. We refer to these uncharged systems as 'neutrachamps'. Although heavy stable charged particles were considered before [1], they were neither put forward as dark matter candidates nor considered within the mass range to which we are led.

To be specific, we make two assumptions that are inessential to the thrust of our arguments: that champs are singly charged, and that the champ population is matter-antimatter symmetric. The remainder of the paper is organized chronologically in the history of the Universe.

In Section 2, we review the arguments of Zel'dovich *et al.* [2] concerning the early Universe annihilation of champs and their consequent relic abundance today. The requirement that $\Omega = 1$ fixes a relation between the champ mass and its annihilation cross-section. Two extreme assumptions about the annihilation mechanism are shown to constrain the champ mass to lie between 2 and 1000 TeV. Our subsequent arguments, and the failure of experiments to find super-heavy isotopes, exclude the lower portion of this window.

The fate of champs during and after the era of nucleosynthesis is considered in Section 3. We elaborate on an argument of Chivukula and Walker [3] to conclude that negative champs emerge from this era bound to protons (as neutrachamps) rather than to α -particles. Section 4 deals with the evolution of primordial fluctuations between the onset of mass dominance and the occurrence of atomic recombination. We show that neutrachamps (but not charged champ states) can emulate collisionless cold dark matter because they are effectively decoupled from thermal fluctuations if their mass exceeds ~ 10 TeV.

Section 5 considers the evolution of galaxies and their halos in the period after recombination. We are led to a lower limit on the champ mass: paradoxically, charged-champ states lighter than ~ 20 TeV cannot remain as a halo, but would fall into the galactic disk. If they are heavy enough, only a few of them should be found within primordial planets or stars. Section 6 is devoted to the search for cosmic champs and the observation of characteristic X -rays signalling the presence of galactic champs, whilst Section 7 considers the search for ambient champs on Earth and in lunar material. The final Section mentions some of the extraordinary sequelæ that would result with the confirmation of our champ hypothesis.

2. Bounds on the Champ Mass from their Relic Abundance

In a champ-antichamp symmetric Universe, conventional Big-Bang theory fixes the present relic champ (and antichamp) abundance, ρ_0 , in terms of the champ mass, M , and its annihilation cross section, σ . We exhibit two extreme (but not necessarily limiting) estimates for σ based on analogies with elementary electrons and with composite protons. The weakly-annihilating champ is taken to be a singly-charged weak singlet without any

strong self-couplings. Its annihilation cross-section is relatively small. For M much greater than the electroweak mass scale, the dominant annihilation mode of the leptonic champ is into two neutral weak intermediaries, for which the cross section, in the non-relativistic limit, is given by:

$$\sigma v = \pi \alpha^2 M^{-2} \cos^{-4} \theta_w, \quad (2.1)$$

where θ_w is the weak mixing angle and v is the relative velocity of the annihilating particles. For our example of rapidly annihilating champs, we simply scale down the measured nucleon-antinucleon low-energy annihilation cross-section by two powers of mass, $\sigma v = (m/M)^2 \sigma_{NN}$, where m is the nucleon mass. Such a relatively large cross section may be expected if champ constituents are endowed with color-like self interactions appropriate to their energy scale. In the velocity range (measured relative to c) $0.05 < v < 0.5$, a rough fit to the data [4] is given by $\sigma_{NN} \simeq (35/v)$ mb, so that the annihilation cross section for baryonic champs, again in the non-relativistic limit, is:

$$\sigma v = (m/M)^2 35 \text{ mb}. \quad (2.2)$$

In both cases, our models involve expressions for σv that are temperature independent in the relevant domain.

The present temperature of the Universe is $T_\gamma = 2.7 \text{ K}$, the Planck mass is $M_{Pl} = 1.2 \times 10^{19} \text{ GeV}$, and the total number of degrees of freedom of the three-family standard model is $g_w = 427/4$ well above the weak scale. Champs ‘freeze out’ at the temperature T^* such that their annihilation rate coincides with the Hubble expansion rate. This occurs [5] at $x \equiv M/T^* \simeq 40$, and results in a prediction of ρ_0 , the present champ (plus antichamp) mass density:

$$\rho_0 = 2x \sqrt{\frac{\pi^3}{45 g_w}} \frac{T_\gamma^3}{M_{Pl} \sigma v}. \quad (2.3)$$

For the Universe to be champ dominated and its mass to be critical, we must have $\rho_0 \simeq \rho_c \simeq 2 \times 10^{-29} h^2 \text{ g cm}^{-3}$, with the Hubble-constant controversy still raging between $h \sim \frac{1}{2}$ and $h \sim 1$. The champ annihilation cross section (times relative velocity) is determined by (2.3) to be $\sigma v = 3 \times 10^{-38} h^{-2} \text{ cm}^2$. It follows from (2.1) that our leptonic champ must have a mass $M \simeq 2 h \text{ TeV}$, and from (2.2) that our baryonic champ must have a mass $M \simeq 1 h \text{ PeV}$. (Should nature display a cosmological champ-antichamp asymmetry, these estimates become mere upper limits.) Champs, if they exist at all, may have intermediate (or even more extreme) properties. The point of this Section is only to show that cosmic

champs, within a wide range of masses, can provide the critical mass density if they have a suitable annihilation cross section. Further cosmological constraints on M , to be discussed presently, provide a lower limit $\sim 20 \text{ TeV}$. Since it is difficult to imagine what mechanism could yield an annihilation cross section larger than our baryonic *ansatz*, we henceforth focus on a range of champ masses extending from 20 TeV to 1000 TeV.

3. The Fate of Champs during Primordial Nucleosynthesis

Positive champs survive the era of nucleosynthesis unscathed. Much later, at recombination, they acquire electrons to form super-heavy isotopes of Hydrogen. The fate of negatively-charged champs is more complex. Their binding energies to nuclei determine the temperature and epoch at which champ-nucleus composites may form. The relevant nuclei in the early Universe are protons (for which the champ binding energy [1] is $E_B \simeq 25 \text{ keV}$), and α -particles (for which the binding energy is $E_B \simeq 311 \text{ keV}$). Since these energies are smaller than nuclear binding energies, champ nucleosynthesis occurs well after conventional nucleosynthesis. Furthermore, since the champ number density is small compared to the baryon number density, the presence of champs cannot significantly affect nucleosynthesis.

Our discussion of champ nucleosynthesis parallels [3], in which constraints are deduced on the nature of hypothetical charged technicolored relics. Let $n_\gamma = 2\zeta(3)T^3/\pi^2$ be the photon number density, $n_B \equiv \eta n_\gamma$ the baryon number density, $Y_N \equiv n_N/n_B$ (with $N = p, \alpha, \dots$) the fraction of baryons in the form of nuclei N of atomic mass A and atomic number Z , n_{Nc} the number density of charmed nuclei, and n_c that of free negative champs. In thermal equilibrium at temperature $T < E_B$, the fraction of champ bound states is dominated by the ground-state contribution and is given by:

$$\frac{n_{Nc}}{n_c} = \sqrt{32/\pi} \zeta(3) \eta Y_N \left(\frac{T}{Am} \right)^{3/2} e^{(E_B/T)}. \quad (3.1)$$

The relic abundance of a particular champ-nucleus bound state is determined by (3.1) evaluated at the temperature at which the rate of the capture process falls behind the expansion rate of the Universe, which is to say, when:

$$n_N(\sigma v) = \sqrt{\frac{4\pi^3 g_*}{45}} \frac{T^2}{M_{Pl}}, \quad (3.2)$$

where $g_s \simeq 3.91$ is the effective number of relativistic species (photons and three kinds of neutrinos), and $\langle \sigma v \rangle$ is the $\mathcal{NC}^-$ radiative capture cross section, both at photon temperature T .

Radiative capture of an atomic nucleus by a negative champ is much like recombination of an electron and a nucleus. The recombination cross section to the champ-nucleus ground state, in a Coulomb model which neglects the nuclear size, may be deduced from the atomic result [6]. The thermally-averaged product of velocity and cross section, in this approximation [3], becomes:

$$\langle \sigma v \rangle \simeq \frac{1024 \pi \sqrt{2\pi}}{3 e^4 A m \sqrt{A m T}} \alpha Z^2 E_B^2 r_B^2, \quad (3.3)$$

where the ground-state binding energy is $E_B = \frac{1}{2} \alpha^2 m A Z^2$, and the ‘Bohr radius’ is $r_B = (\alpha m A Z)^{-1}$. We stress that this approximation, while it is appropriate to describe the capture of a proton by a champ, is certainly inadequate to describe correctly the capture of an α -particle. In the latter case, the nuclear size is about half of the ‘Bohr radius’ of the champ-nucleus bound state and cannot be neglected. Moreover, the binding energy, computed with a distributed charge density [1], is ~ 311 keV, whilst its Coulomb value is $\sim 8\alpha^2 m \simeq 400$ keV.

Plunging onwards nonetheless, we insert (3.3) into (3.2) to deduce a preliminary estimate of the freeze-out temperature:

$$T_f \simeq 0.31 \frac{g_s A^3 m^3}{M_{Pl}^2 Y_{\mathcal{N}}^2 \eta^2 Z^8 \alpha^6}. \quad (3.4)$$

Eq. (3.4), evaluated for α ’s and for protons, yields:

$$T_f(\alpha) \simeq 7.7 \text{ keV} (6 \times 10^{-10} / \eta)^2, \quad (3.5)$$

$$T_f(p) \simeq 137 \text{ eV} (6 \times 10^{-10} / \eta)^2, \quad (3.6)$$

where we have put $Y_\alpha \simeq 1/16$ and $Y_p \simeq 15/16$. Alpha freeze-out occurs first. Eq. (3.1), evaluated at $T_f(\alpha)$, depends sensitively [3] on the baryon abundance. For $\eta > 5.8 \times 10^{-10}$, and using our tentative estimate of $T_f(\alpha)$, we find that most champs bind to α -particles. For smaller values of η , most champs avoid α -capture. However, they do not survive as bare champs since (3.1) at $T_f(p)$ guarantees their subsequent capture by protons. Depending

on the value of η , essentially all negative champs bind to protons (forming neutrachamps) or to α ’s (forming positively-charged ‘isotopes’ of $\mathcal{C}^{++}$).

The lamentable uncertainty in the fate of negative champs is an artifact of our use of the Coulomb approximation, which accidentally leads to a borderline value for the recombination cross section of champs and α -particles. We believe that a correct estimate of α -champ recombination yields a significantly smaller cross section than the Coulomb value, and hence a higher value of $T_f(\alpha)$. In lieu of a precise computation of the recombination coefficient for a distributed charge density, we offer an intuitive argument. We ‘correct’ the Coulomb cross section (3.3) by replacing the phase-space factor E_B^2 by $(311 \text{ keV})^2$. This is its value calculated [1] with a potential taking into account the extended charge distribution of the α -particle. We replace the geometric factor r_B^2 in (3.3) by its value in a Coulomb potential that has been reduced in strength so as to reproduce this calculated binding energy. The net result is to decrease the recombination cross section by the factor $\sim (311/400)$, and thus to increase $T_f(\alpha)$ by $\sim 65\%$. This value of the freeze-out temperature ensures that negative champs do not bind to α -particles. They remain bare until the time is ripe for them to bind to protons. If our argument is valid, virtually all negative champs emerge from the era of champ-nucleosynthesis in the form of neutrachamps.

Hardly any champs bind to larger nuclei whose primordial abundances are far too low for equilibrium with champs to set in. At later times, some champs may migrate into hot stars where they may be reprocessed. Neutrachamps, and the probably quite rare $\alpha \mathcal{C}^-$ systems, could dissociate, with the champs finding their way to high- Z nuclei, and eventually, to the primordial Earth. However, we show in Section 5 that only a few charged champs, and hardly any of the neutrachamps, can collapse into the galactic disk if they are sufficiently heavy. In Section 6, we focus on the champ abundance on Earth due to the ‘known’ influx of cosmic champs, should they comprise the non-baryonic dark matter of the Milky Way.

4. The Growth of Fluctuations before Recombination

If $\Omega = 1$, the onset of mass dominance takes place well before recombination. During the extensive period between these two transitions, primordial density fluctuations

¹ Recent analyses [7] of nuclear abundances, especially of ^7Li , suggest a value of η smaller than 4×10^{-10} . This result would favor the likelihood that almost all negative champs emerge as $p \mathcal{C}^-$ bound states, neutrachamps.

can cease being adiabatic, and collisionless cold dark matter can develop relatively large fluctuations without leaving any thermal imprint. Thereby can the absence of observed anisotropies of the cosmic black-body radiation be reconciled with the (suspected) appearance of galaxies at $z \sim 4$. The question we must address is whether champs can do the same thing even though they are charged and hardly collisionless.

The origins of galaxies and larger structures, their peculiar motions, and the high degree of uniformity of the cosmic background radiation are not fully understood. A scenario [8] that fares relatively well in explaining these observations is that of a critical inflationary Universe with its present mass density dominated by collisionless cold dark matter. Inflation presumably initiates galaxy formation by furnishing adiabatic density perturbations at a very early time. Once matter dominance occurs, these primordial fluctuations may depart from adiabaticity. The dark-matter components of over-dense fluctuations may grow under the effect of their own gravity to become the seeds of galactic halos. On the other hand, the radiation and baryonic components of fluctuations of galactic size are washed away by Silk [9] damping by the time of recombination, and leave no large thermal imprint. Only at a later stage does the baryonic density contrast catch up with that of the dark matter (as baryons fall into the dark potential wells) to end up as the visible component of galaxies.

The dissipationless nature of dark matter is crucial. But, champs interact electromagnetically with photons and with charged particles in the primordial plasma: they are neither collisionless nor dissipationless. Champ-dominated density fluctuations face the same threats as those of ordinary matter: to be washed away by a mismatch between radiation pressure and matter density (Silk damping), or to leave a falsifiable signature in the form of non-uniformities of the background radiation. Sufficiently heavy champs, however, are not efficiently dragged around by ordinary matter and radiation. We now estimate the minimum mass that singly-charged champs (either C^+ , or, αC^-) must have to qualify as dissipationless dark matter. Subsequently, we consider the rôle that neutrachamps play in galaxy formation, since we have seen in Section 3 that this is the likely form in which negative champs emerge from the era of champ nucleosynthesis.

When matter dominance begins, the era of ‘stagnation’ ends and champ fluctuations can begin to grow. Assume, contrary to our expectations, that all negative champs are bound to α -particles. The plasma, then, has four ingredients: positive champ states (C^{++} ’s and αC^{++} ’s), baryons, electrons and photons. Adiabatic fluctuations of galactic size, in which all components move together, are suppressed by Silk damping, while the strength

of the Coulomb force guarantees that no charge separation can take place. We are left with two possible non-adiabatic modes: those in which champs disengage from baryons (with electrons establishing charge neutrality), and those in which champs, baryons, and electrons move together relative to the thermal background.

The former mode is damped by Coulomb collisions between champs and protons. Spitzer [10] computes the characteristic equilibration time for this process to be:

$$t_1 = \frac{3 M m}{8 \sqrt{2\pi} \alpha^2 n \ln(3T/\alpha k_D)} \left[\frac{T}{m} + \frac{T_c}{M} \right]^{3/2}, \quad (4.1)$$

where M is the champ mass, m is the proton mass, n is the number density of protons, and $k_D = (4\pi n_e \alpha / T)^{1/2}$ is the Debye momentum at this density and temperature. The equilibration time for champ fluctuations is:

$$t_1 \simeq 10^7 (M/\text{PeV}) (\text{eV}/T)^{3/2} \text{ s}. \quad (4.2)$$

At recombination, $T_n \simeq 0.25 \text{ eV}$ and the age of the Universe is $t_n \sim 10^{13} \text{ h}^{-1} \text{ s}$, an appropriate dynamical time scale. Survival of champ-baryon fluctuations requires that $t_1 > t_n$ at T_n , which, according to (4.2) obliges $M > 10^5 \text{ PeV}$. This result is far beyond the upper limit set in Section 2. We conclude that fluctuations in which charged champs separate from baryons are too rapidly damped to play any significant rôle in galaxy formation.

Undulations in which champs, baryons and electrons move together relative to photons are damped by the Compton scattering of photons off the charged constituents of the plasma. The dominant effect is due to scattering off electrons. Using a computation of Compton drag due to J.I. Katz [11], we obtain for the equilibration time of these fluctuations:

$$t_2 = \frac{135 \rho m_e^2}{64 \pi^3 \alpha^2 T^4 n_e}, \quad (4.3)$$

where ρ is the fluid density (dominated by champs), m_e is the electron mass, and n_e is its number density. For the case at hand, the equilibration time becomes:

$$t_2 \simeq 7 \times 10^9 (\text{eV}/T)^4 (0.03/\Omega_B) \text{ s}, \quad (4.4)$$

where Ω_B is the fraction of ρ_c in baryons. Our result is independent of the champ mass since it depends only on the bulk density of the fluid, which is the critical mass density at temperature T . Again, the equilibration time t_2 is short compared to t_n at T_n . This type of fluctuation is also damped out and useless for galaxy formation.

We are driven to conclude that a Universe dominated by charged champ states is no better (nor worse) than one dominated by baryons with respect to the formation of proto-galaxies prior to recombination. However, negative champs are apt to be in the form of the more weakly interacting neutrachamps, which, we now show, can serve as conventional collisionless cold dark matter.

The dominant mechanism for a neutrachamp to adapt to the local motion of the primordial fluid is via its collisions with protons and α -particles, the only abundant nuclei. The estimate of dE/dx for these processes, discussed in detail in Section 6, is a difficult art. We cannot use the Born approximation result, (6.10), since the collision velocities are merely $\sim 10^{-5}$. A result of Lindhard *et al.* [12] applicable to the stopping power of slow ions may be modified to describe neutrachamp collisions with nuclei to yield (6.6). Thus, the energy loss of a slow neutrachamp is expected to be:

$$\left. \frac{dE}{dx} \right|_{\mathcal{N}} = \frac{\pi^2 Z \rho_{\mathcal{N}}}{A e m^2}, \quad (4.5)$$

and the corresponding equilibration time is:

$$t_{\mathcal{N}} \equiv \frac{dt}{d \ln E} = M \left(\frac{2T}{\pi A m} \right)^{1/2} \frac{dx}{dE} = \frac{M e \sqrt{2} T A \bar{m}^3}{\pi^{3/2} Z Y_{\mathcal{N}} \rho_{\mathcal{N}} \Omega_b} \left(\frac{T_{\gamma}}{T} \right)^3, \quad (4.6)$$

where the result is put in terms of the present critical mass density ρ_c and the photon temperature T_{γ} . Protons dominate over α 's in (4.5). Eq. (4.6), evaluated at the recombination temperature $\simeq 0.25$ eV, becomes:

$$t_{\mathcal{N}} = \left(\frac{M}{10 \text{ TeV}} \right) \left(\frac{0.01}{\Omega_b h^2} \right) 2.0 \times 10^{13} \text{ s}. \quad (4.7)$$

For $M > 10$ TeV and for acceptable values of $\Omega_b h^2$, the neutrachamp equilibration time exceeds the age of the Universe at recombination.

At the onset of mass dominance, positive champs and neutrachamps, in equal numbers, provide the critical mass density. We have shown that neutrachamps with plausible masses are sufficiently noninteracting to act as dissipationless cold dark matter. Their fluctuations, just like those of collisionless cold dark matter, grow prior to recombination. Positive champs, up to this time, remain as uniformly distributed as baryonic matter and photons, as observations of the homogeneity of the cosmic black-body radiation indicate.

5. Bounds on the Champ Mass from Galactogenesis

Baryonic matter lying within identifiable stars, dust and gas constitutes only a few percent of critical mass density. Large spiral galaxies, like the Milky Way, have non-luminous halos of mass and size at least 5 to 10 times larger than those of their visible and manifestly baryonic components. Inflationary solutions to various cosmological problems imply that the present mean energy density, ρ , is very close to the critical mass density: $\Omega \equiv \rho/\rho_c \simeq 1$. Successful predictions of the primordial abundances of light elements seem to require that the baryonic contribution to Ω be bounded by $\Omega_b < 0.2$. The contribution of galaxies to Ω , with their dark halos included, is estimated to be as large as $\Omega_o \sim 0.2$. Although galactic halos conceivably may be made of ordinary matter, we explore in this and subsequent sections the consequences of assuming that champs constitute their unseen components.

The cold dark matter scenario for the late stages of the genesis of galaxies is both plausible and well developed. We review it to ascertain whether champs (as the dominant component of galactic halos) may serve as well as collisionless cold dark matter. After recombination, a galaxy is an over-dense concentration of dark matter that lags behind the universal expansion rate due to the effect of its own excessive gravity. It eventually recollapses onto itself, while ordinary matter (whose distribution on galactic scales has been made uniform by Silk damping) falls into the dark matter potential wells soon after recombination, leaving no discernible inhomogeneity in the microwave background. The galaxy, when it begins to recontract, is a cold and uniform mixture of dark matter and of whatever Hydrogen and Helium have fallen within.

Gravitational infall stops when the system reaches about half its original size by the process of 'violent relaxation', whereby velocities are redistributed as each particle travels through the rapidly-varying gravitational fields of its neighbors. The resulting distribution is approximately Maxwellian, except that the velocity dispersion of all particles is the same, independent of their mass. (Such an entity is often dubbed an 'isothermal sphere', although that is just what it is not.) This description is correct only insofar as forces other than gravity can be neglected, which is not the case for the ordinary-matter component of the galaxy. In the proto-galactic crunch, the temperature becomes high enough to ionize ordinary matter, which radiates efficiently and consequently collapses into a disk or bulge that becomes a luminous galaxy. Collisionless dark matter neither radiates nor otherwise loses energy. It remains as an extended spherical halo, with a flat rotation

curve characteristic of an isothermal polytrope, and with a characteristic virial velocity, $v \sim 200\text{--}300 \text{ km s}^{-1}$. This is also the typical velocity of the protons in the ionized gas, corresponding to a temperature $T \sim 10^6 \text{ K}$. For this scenario to succeed, the cooling time of the plasma must be shorter than the dynamical time, the ‘year’ of a test particle in galactic orbit. The upper limit to a galaxy mass corresponds to approximate equality between the two relevant time scales. Our own galaxy is not far from this limit. Its mass and radius, including the halo, are $M_G \sim 10^{12} M_\odot$, $R_G \sim 100 \text{ kpc}$.

In the previous Section we saw that neutrachamps, prior to recombination, snugly fit into the conventional cold dark matter scenario for the growth of the density fluctuations that become the seeds of galactic halos. We now show that this successful imitation extends to the later stages of galactic evolution.

At recombination, fluctuations in neutrachamp density of galactic size must have grown to a contrast $\delta_0 \equiv \delta\rho/\rho \sim 10^{-2}$, whilst fluctuations in C^+ ’s, ordinary matter, and photon temperature, since they have not been seen in the cosmic microwave background, must be at least two or three powers of ten smaller [8]. Let δ_+ and δ_n be the density contrasts in C^+e^- (super-heavy hydrogen atoms) and in ordinary atoms. Let Ω_0 , Ω_+ , and Ω_n be the fractions of critical density in neutrachamps, heavy Hydrogen, and baryons, respectively. (In a critical density Universe, such as we assume, their sum is unity.) The equations governing the evolution of the different δ ’s in the linear régime are:

$$\delta_i'' + \frac{3}{2\mathcal{R}} \delta_i' = \frac{3}{2\mathcal{R}^2} [\Omega_0 \delta_0 + \Omega_+ \delta_+ + \Omega_n \delta_n] \quad (i = 0, +, n), \quad (5.1)$$

where primes denote derivatives relative to $\mathcal{R}$, the scale size of the Universe. The solution of (5.1) is:

$$\delta_i = A \mathcal{R} + B \mathcal{R}^{-3/2} + a_i + b_i \mathcal{R}^{-1/2}, \quad (5.2)$$

where

$$\sum \Omega_j a_j = \sum \Omega_j b_j = 0. \quad (5.3)$$

Let the initial conditions at recombination, $\mathcal{R} = \mathcal{R}_n \sim 1/1400$, be $\delta_0 = \delta_0^n \sim 2 \times 10^{-2}$, $\delta_+^n = \delta_n^n \sim 0$, and let the Universe be critical and champ-antichamp symmetric: $\Omega_+ = \Omega_0 = (1 - \Omega_n)/2$. Eq. (5.1) has a non-decreasing solution satisfying these boundary conditions with coefficients $B = b_i = 0$, and

$$\begin{aligned} A &= (1 - \Omega_n) \delta_0^n / (2\mathcal{R}_n) \\ a_0 &= (1 + \Omega_n) \delta_0^n / 2 \\ a_+ &= a_n = -(1 - \Omega_n) \delta_0^n / 2. \end{aligned} \quad (5.4)$$

This result establishes that the density contrasts in super-heavy hydrogen and in ordinary matter have practically caught up with the initial neutrachamp density contrast ($\delta_+ = \delta_B \sim \delta_0$) by the time $\mathcal{R} = 10 \mathcal{R}_n$. Furthermore [8], galaxies ‘go nonlinear’ ($\delta = \delta_c \sim 4.6$) when $1 + z = \mathcal{R}^{-1} \sim \delta_0^R / 2\delta_c \mathcal{R}_n \sim 3$. This script closely follows the standard dark matter scenario in which only baryonic matter falls into the dark traps. (The only difference in our analysis is that δ_0^n must start off at twice its conventional value in order for galaxies to reach the nonlinear regime at a similar redshift.)

In the course of violent relaxation, the proto-galaxy becomes hot enough to ionize atoms. In this environment, charged champs suffer dissipative interactions akin to those of protons and electrons. The charged-champ component of the halo (though certainly not its more weakly interacting neutrachamp component) incurs the danger of being dragged down with ordinary matter into the visible disk of the galaxy. Sufficiently heavy champs would be stiff enough to escape this menace. How massive must they be? The most stringent limit stems from the existence of halos of large galaxies whose ordinary matter remains ionized, exerting frictional drag on charged champs for the longest time. We must compare the characteristic time for a champ to lose its kinetic energy in a plasma to a dynamical time, the orbital period of a champ. (We refer to one and not several orbital periods because the cooling time of a large galaxy is comparable to its dynamical time and the champ encounters only neutral matter on its second coming.)

Consider a large galaxy, such as ours, with a baryonic mass $\sim 10^{11} M_\odot$. At proto-galactic collapse, ionized gas may have extended over a region of radius $R \sim 100 \text{ kpc}$, corresponding to an average number density $n_e \sim n_p \sim 10^{-3} \text{ cm}^{-3}$. The dynamical time, $t_{dyn} = (G\rho)^{-1/2}$, for such a galaxy is $\sim 10^9 \text{ y}$, where we assume a 10 to 1 ratio of dark to ordinary matter. During this time, a charged champ may lose energy by making Coulomb collisions with electrons or nuclei. The characteristic time t_{mat} for a champ of velocity v to lose energy in such collisions may be estimated with (4.1). The result is dominated by collisions with protons. Putting $T \sim 10^6 \text{ K}$, and $T_c \sim \frac{1}{3} M v^2$ with $v \sim 10^{-3}$, we obtain $t_{mat} \sim 5 \times 10^8 (M/10 \text{ TeV}) \text{ y}$. The condition $t_{mat} > t_{dyn}$ restricts champs to be heavier than $\sim 20 \text{ TeV}$ if C^+ ’s are to remain up in the halo and not to collapse.

The analog of (4.3) for the characteristic time for charged champs to lose energy due to Compton scattering [11] off the photons of the crunched galactic plasma is

$$t_\gamma = \frac{135 M^3}{64 \pi^3 \alpha^2 T^4} \quad (5.5)$$

For $T = 10^6 K$, the result is $t_\gamma = 5 \times 10^{11} (M/10 \text{ TeV})^3 \text{ y}$. The effect of Compton drag on champs is negligible.

Given the rough character of our considerations and the limitations of our understanding of galaxy gestation, the limit $M > 20 \text{ TeV}$ for the survival of charged champs in dark halos (and the $M > 10 \text{ TeV}$ limit deduced in Section 4 for the survival of neutrachamp galactic seeds) must both be taken *cum* the usual astrophysical *grau* *salis*. These limits contradict the upper limit $M < 2 h \text{ TeV}$ for the mass of leptonic champs obtained in Section 2. The existence of such particles seems to be totally excluded, while the existence of more strongly self-coupled champs is a viable possibility in the range of masses below 1 h PeV and above a shakier 20 TeV : a window for terrestrial experiments to close or to fill.

At the onset of violent relaxation (the process that segregates baryonic matter from the champ halo), the halo consists of its aboriginal neutrachamps together with a comparable number of C^+ 's that are swept in with the baryons. While there is no danger that the relatively weakly interacting neutrachamps collapse into the disk, the constraint $M > 20 \text{ TeV}$ is needed to ensure that most C^+ 's remain suspended in the halo, segregated from baryonic matter, and consequently, rare on Earth and in the stars. (Inevitably, some C^+ 's do collapse into the galactic disk, perhaps becoming a significant stellar component.) We conclude our discussion of galaxy formation on an upbeat note: The champ hypothesis is viable and champs can be just as effective as collisionless cold dark matter to explain the origins of large scale structure.

6. To Search for Cosmic CHAMPS

The local density of non-baryonic dark matter in the solar neighborhood is thought to be²

$$\rho_{\sim\odot} \sim 4 \times 10^{-25} \text{ g cm}^{-3}. \quad (6.1)$$

We assume that this material consists of champs with typical virial velocities about the galaxy. With velocities relative to Earth $v \sim 10^{-3}$, their flux upon the surface must be

$$F \sim 2 (\text{PeV}/M) \text{ cm}^{-2} \text{ s}^{-1}. \quad (6.2)$$

² Published estimates [13] of the non-baryonic dark matter density in the solar neighborhood lie in the range $3 - 7 \times 10^{-25} \text{ g cm}^{-3}$. We use a conservative intermediate value.

This extraordinarily large flux ought to be easily detectable, or so it may seem. There are three ways to find champs: by the direct detection of cosmic champs via their interactions in flight, by the detection of a characteristic electromagnetic signal produced by cosmic champs, or by the search for ambient champs mimicking super-heavy isotopes in terrestrial, lunar, or meteoritic materials. In this Section, we address the former challenges. To devise a suitable search protocol, the interactions of cosmic champs with ordinary matter must be understood.

What is the fate of positive champ states (C^+ and αC^-) that enter the atmosphere? Because the binding energy of αC^- is so large ($\sim 311 \text{ keV}$), it cannot dissociate at typical collision energies, and because of its positive charge there is little chance that the C^- may be exchanged to bind to a higher- Z nucleus. Therefore, positive champ states interact and lose energy exclusively by Coulomb collisions with nuclei and electrons.

A computation of the energy loss of slowly-moving massive charged particles was carried out [12] by Lindhard *et al.* Their result is in reasonable agreement with experiment [14] down to ion velocities $\sim 10^{-3}$, and is easily adapted to the case of an exceedingly heavy incident charged champ to yield:

$$\frac{dE}{\rho dx} \sim \frac{\pi}{A m m_e} \left[\frac{8v}{\alpha} + \frac{\pi}{e} Z^{2/3} \right]. \quad (6.3)$$

The first bracketed term corresponds to electronic collisions, the second to nuclear collisions. For $v = 10^{-3}$, and on a Nitrogen target, (6.3) yields an energy loss of $\sim 525 \text{ MeV cm}^2 \text{ g}^{-1}$, mostly due to nuclear collisions. This result agrees with that obtained in Born approximation [15]. At such a speed, cosmic champ with a mass at the upper limit of 1000 TeV has a kinetic energy of $\sim 500 \text{ MeV}$. Charged champs are stopped by a mere fraction of a gram cm^{-2} . They are thermalized in the top thousandth of Earth's atmosphere and their detection in flight would require the deployment of unshielded detectors mounted on satellites or high-flying rockets³.

Some positively-charged cosmic champs may be ionized within their intragalactic environment, and some of these may be subject to the accelerative mechanism that produce

³ A high-altitude search for neutral halo particles with typical strong interaction cross sections using a Silicon detector has been carried out [16]. It yields a limit $M > 100 \text{ TeV}$ for an assumed isotropic nuclear cross section $\sigma \sim 0.5 \text{ barn}$. This limit is not applicable to charged champs because they would be stopped by the 4.5 g cm^{-2} of air remaining above the rocket at its maximum altitude.

energetic cosmic rays. Since champs are few compared to protons, and since only a few galactic protons are in the form of cosmic rays, we expect that only a tiny fraction of champs acquire sufficient energy to penetrate Earth's atmosphere in flight.

In Section 3 we argued that negative champs emerge from the early Universe bound to protons. The behavior of these neutrachamps is more complex but less well understood than that of their charged brethren. They may engage in the following interactions:

The neutrachamp is a relatively large system consisting of a central point-like negative charge surrounded by an orbiting proton with a 'Bohr radius' $(m\alpha)^{-1} \sim 30$ fm. We consider first the energy loss of slow neutrachamps due to elastic collisions as they pass through matter. We offer three distinct estimates.

We find a rough lower limit to its energy loss by treating the neutrachamp as if it were a rigid scattering center with a shielded Coulomb potential. In the center-of-mass system, the target nucleus scatters [17] approximately backwards with a geometric cross section $\sim \pi(\alpha m)^{-2}$. This mechanism provides a small energy loss for neutrachamps:

$$\frac{dE}{\rho dx} \simeq \frac{2\pi v^2}{\alpha^2 m^2} \simeq 0.03 \text{ MeV } (10^3 v)^2 \text{ cm}^2 \text{ g}^{-1}. \tag{6.5}$$

However, the pC^- system is not rigid, but may be polarized during its slow encounter with the target nucleus at impact parameters larger than $(\alpha m)^{-1}$. A more plausible estimate of energy loss which takes this effect into account is obtained by following an analogy with Lindhard's analysis [12] addressing the scattering of slow ions by neutral atoms in a classical approximation with an r^{-2} potential. In our case, the proton shielding the champ plays the rôle of the electrons shielding a nucleus. Transcribing the nuclear component of (6.3) to the case at hand, we find:

$$\frac{dE}{\rho dx} \simeq \frac{\pi^2 Z}{A \epsilon m^2} \simeq 0.46 (2Z/A) \text{ MeV cm}^2 \text{ g}^{-1}. \tag{6.6}$$

A third estimate, following Mott and Massey [15], should be valid for larger values of A and v . In Born approximation, we write:

$$\frac{d\sigma}{d\Omega} = \frac{A^2 m^2}{4\pi^2} \left| \iint d^3x d^3y e^{i\mathbf{q}\cdot\mathbf{x}} Z\alpha \left(\frac{1}{|\mathbf{x}|} - \frac{1}{|\mathbf{x}-\mathbf{y}|} \right) \phi^2(\mathbf{y}) \right|^2, \tag{6.7}$$

where ϕ is the Coulomb wave function of the bound proton in its ground state, and $\mathbf{q} = \mathbf{k}' - \mathbf{k}$ is the momentum transfer to the neutrachamp. Using (6.7), we compute for the energy loss:

$$\frac{dE}{\rho dx} = \frac{2\pi Z^2}{m^2} \mathcal{F}(A^2 v^2 / \alpha^2), \tag{6.8}$$

with

$$\mathcal{F}(s) = \frac{1}{s} \int_0^s \frac{p(2+p)^2}{(1+p)^4} dp. \tag{6.9}$$

For collisions with Nitrogen atoms at $v = 10^{-3}$, we find $\mathcal{F} \simeq 0.41$, and an energy loss far exceeding the previous estimate:

$$\frac{dE}{\rho dx} = 32 \text{ MeV cm}^2 \text{ g}^{-1}. \tag{6.10}$$

To determine whether any neutrachamps reach Earth's surface at speed, we must know which of our three estimates is applicable, (6.5), (6.6), or (6.10), since the range of a PeV neutrachamp with incident speed $v \simeq 10^{-3}$ is about equal to the depth of the atmosphere according to (6.6), but is 15 times greater according to (6.5), and 70 times smaller according to (6.10). For the Born approximation to be applicable, the de Broglie wavelength of the nucleus must be far smaller than the size of the neutrachamp, or $Av > \alpha$. In our case, the numbers are comparable, and the Born approximation probably yields an overestimate. Consequently, we favor a result intermediate between (6.6) and (6.10), which suggests that cosmic-ray neutrachamps should be searched for at high altitudes or in space. This view is strengthened by the following discussion of inelastic processes that allow the transmutation of neutrachamps into charged champ states during their atmospheric transit.

The exchange reaction, in which the C^- is transferred from a proton to a larger nucleus, is exothermic. For example, if the target nucleus is ^{14}N , the energy release (i.e., the difference of the binding energies [1] of the champ to Nitrogen and to a proton) is $Q \simeq 3.5$ MeV. The resulting highly-charged $\mathcal{N}C^-$ system is quickly brought to rest by (6.3), amplified by a factor $\sim Z$. We estimate the cross section for exchange in the center-of-mass frame in which the champ is essentially at rest. Let the initial state wave function be $\Psi_i = e^{i\mathbf{k}\cdot\mathbf{x}} \phi_p(\mathbf{y})$, where $\mathbf{x}$ is the position of the incident nucleus and $\mathbf{y}$ that of the

bound proton. We take for ϕ_p a ground-state Coulomb wave function. Let the final state wave function be $\Psi_f = e^{i\mathbf{q}\cdot\mathbf{r}} \phi_{\mathbf{q}}(x)$. It is appropriate [1] to use a harmonic oscillator wave function for a champ bound to a nucleus as large as ^{14}N . Although we use a plane wave state for the final proton, we include a Gamow factor $G(q)$ to describe its final state interaction. The matrix element $\mathcal{M}$, in Born approximation [15], is

$$\mathcal{M} = \int \int d^3x d^3y \Psi_f^* \frac{Z\alpha}{|\mathbf{x} - \mathbf{y}|} \Psi_i \quad (6.11)$$

We obtain, for the exchange cross section in this approximation:

$$\begin{aligned} \sigma_{\text{exchange}} &\sim \frac{1}{\pi} \frac{q}{k} m^2 A |\mathcal{M}|^2 G(q) \\ &\sim v^{-1} \frac{10^{1/4}}{2} \pi^{1/2} \frac{3^{1/2}}{\alpha} \frac{17^{1/4}}{(m\tau_0)^{9/4}} Z^{8/4} \frac{G(q)}{m^{1/2} Q^{3/2}}, \end{aligned} \quad (6.12)$$

where $G = z/(e^2 - 1)$, with $z = 2\pi(Z - 1)\alpha/v_p$ and $v_p \simeq (2Q/m)^{1/2}$. The nuclear radius is $\simeq \tau_0 A^{1/3}$ with $\tau_0 \simeq 1.2$ fm. The resulting reaction cross section off Nitrogen is $\sigma \sim 1.6 \times 10^{-28}$ cm². This is a large result, but not quite sufficient to entirely eliminate the sea-level flux of neutral champs. (Of course, as they come to rest in matter, neutrachamps are surely subject to exchange.)

The remaining interactions in (6.4) are endothermic. Dissociation requires $Av^2 > \alpha^2$, while excitation to the first excited state requires $Av^2 > 0.75\alpha^2$. Consider a collision for which the dissociation reaction is kinematically allowed. Although we might anticipate a geometric cross section, $\sigma_{\text{dis}} \sim \pi(\alpha m)^{-2}$, or some tens of barns, let us carry out a more precise calculation in Born approximation. In the rest system of the pC^- , let k be the momentum of an incoming nucleus, p its final momentum, and q that of the dissociated proton. An approximate expression for the dissociation cross section is:

$$\sigma_{\text{dis}} = \frac{2\pi}{k} m A \frac{1}{(2\pi)^5} \int \int d^3p d^3q |\mathcal{M}|^2 \delta\left(\frac{q^2}{2m} + \frac{p^2}{2mA} - Q\right), \quad (6.13)$$

where $Q = \frac{1}{2}mA(v^2 - v_0^2)$ is the available energy, with $v_0 = \alpha A^{-1/2}$ the threshold velocity for dissociation. The matrix element, $\mathcal{M}$, is given by:

$$\mathcal{M} = \int \int d^3x d^3y e^{-i\mathbf{k}\cdot\mathbf{x}} \phi(y) \frac{Z\alpha}{|\mathbf{x} - \mathbf{y}|} e^{i\mathbf{p}\cdot\mathbf{x} + i\mathbf{q}\cdot\mathbf{y}}. \quad (6.14)$$

The result for the dissociation cross section is:

$$\sigma_{\text{dis}} = \frac{2^6 \pi Z^2}{\alpha m^2 v A^{1/2}} \left(\frac{v}{v_0} + \frac{v_0}{v} - 2 \right) \Theta(v - v_0). \quad (6.15)$$

This is an enormous cross section, amounting to thousands of barns at speeds 10% over v_0 . Air contains $\sim 1\%$ Argon for which $v_0 \simeq 1.15 \times 10^{-3}$. A significant fraction of neutrachamps (those with at least this speed⁴) dissociate in the atmosphere to become bare C^- 's. Thereafter, they lose energy rapidly via (6.3) and are soon captured by Nitrogen or Oxygen nuclei.

While their flux on Earth is considerable, their properties are such as to make the detection of champs in flight unlikely and nearly impossible in ground-based or underground experiments. Charged champs come to rest (6.3) after traversing a tiny fraction of the atmosphere. Neutrachamps lose energy by (6.5) (6.6), or (6.10), but they also may convert to charged states via dissociation (6.15), or by exchange scattering (6.12). Few, if any, champs will reach the ground at speed. However, the neutrachamp, which consists of a massive central C^- surrounded by a proton in a Bohr orbit, is a kind of new chemical element with its own characteristic spectrum. For example, its Lyman α analog line is at $3\alpha^2 m/8 \simeq 18.7$ keV, while its Balmer α line is at $5\alpha^2 m/72 \simeq 3.5$ keV. As it was long ago in the case of Helium, evidence for the existence of neutrachamps may emerge from astronomical spectroscopy⁵. Cosmic neutrachamps, incident on matter, suffer numerous collisions by which they may be excited. Subsequent de-excitation can produce a characteristic monochromatic X-ray signal.

To estimate the cross section for the excitation reaction, we use (6.7) with the following changes. We replace ϕ^2 with the product of the ground-state and excited-state wave functions. For specificity, we focus on excitation to the $2S$ level. We must also include a factor of p/k , where k is the initial momentum of the nucleus in the neutrachamp rest frame, and p is its final momentum. We define $v_0 = \alpha(3/4A)^{1/2}$ to be the threshold velocity below which the transition cannot take place, and $\lambda = p/k = \sqrt{1 - (v_0/v)^2}$ to be a useful parameter. The excitation cross section in Born approximation is:

$$\sigma_{\text{excit}} = 2^9 \frac{3^{-10} \pi Z^2 A(1 - \lambda^2)}{5 \alpha^2 m^2} \left\{ \left[\frac{A(1 - \lambda)}{3(1 + \lambda)} \right]^{-5} + [\lambda \rightarrow -\lambda] \right\}. \quad (6.16)$$

A typical evaluation of (6.16) demonstrates that excitation cross sections are large, so that each neutrachamp, as it slows down, may be excited and de-excited many times producing

⁴ Neutrachamp velocities relative to Earth cannot exceed the Sun's peculiar velocity plus escape velocity from the Galaxy, or a total of $\simeq 2 \times 10^{-3}$.

⁵ We thank David Eichler for suggesting this possibility to us.

of this material (at least $\frac{1}{3}$ of Earth's accumulation of cosmic champs) should now be in the chemical form of super-heavy water.

Many neutrachamps, perhaps most of them, lose their energy or dissociate in the atmosphere. Consequently, they bind to atmospheric atoms forming super-heavy Carbon and Nitrogen. Those that survive as neutrachamps to reach the oceans are mostly transformed into superheavy Carbon. Those that strike land (or the Moon) are sure to exchange with common elements. They form super-heavy Al, Mg, etc., with abundances proportional to those of elements with Z one greater. The most promising chemical sites for future searches for ambient champs would seem to be in oxidized forms of Hydrogen, Carbon, and Nitrogen.

The continental crust, though subject to geological processing and redistribution, is not thought to have been significantly recycled with mantle material: all of its champ accumulation should still lie somewhere within. Since the mean depth of the continental crust [18] is $\sim 7 \times 10^6 \text{ g cm}^{-2}$, champs should comprise $\sim 10^{-7}$ of its mass, of which a considerable portion should be in the form of super-heavy water, carbonate, or nitrate, or as super-heavy N_2 . We suspect that super-heavy elements are biologically rejected, so that rocks of organic origin are unpromising sites, as are rocks of great age or of magmatic origin. Champ matter on land should be more abundant than gold though less than silver.

The concentration of champs in the oceanic crust is considerably smaller, since this material has a relatively short residency before being subducted into the mantle. However, a slowly accumulating marine sediment may be enriched in champs. Certain Pacific sediments formed $\sim 60 \text{ My}$ ago at the rate of about a millimeter per millenium [19]. We estimate their super-heavy water content, assuming equilibrium between incidence and sedimentation (and that the deposit is not leached out) to be $\sim 10^{-7}$ by mass.

Meteorites could be a promising material in which to find champs. However, ablation on re-entry may effectively purge them of what surface deposition they originally had. Lunar material should be the richest of all champ ores. The Moon [18] has been geologically quiet for $\sim 3.8 \text{ Gy}$, during which time its surface was distressed only by meteoritic impacts, and to an average depth no greater than 10 m. All of the incident champs remain concentrated near the lunar surface, so that the champ content of lunar material should be $\sim 10^{-4}$ by mass.

There have been many searches for super-heavy isotopes. We ardently advocate their continuation and extension to larger masses than have been considered to date. While our champ model is put forward to be champ-antichamp symmetric, this assumption is

many X -rays. The excitation cross section to the $2S$ level of Oxygen is ~ 40 barns at $v = 1.1 v_0 \simeq 1.74 \times 10^{-3}$, and ~ 400 barns at $v = 1.5 v_0$. (This is yet one more mechanism that can contribute to the energy loss of cosmic neutrachamps.)

Consider cosmic neutrachamps incident upon Earth, Moon or Sun. The threshold velocity for neutrachamp excitation off Oxygen is $v_0 \simeq 1.58 \times 10^{-3}$. Many neutrachamps are expected to exceed this velocity relative to the Earth or the Moon, and all those incident upon the Sun surely do, for its escape velocity is $\sim 2 \times 10^{-3}$. Slower neutrachamps can produce other characteristic X -rays via the exchange-excitation reaction (6.4), or via dissociation followed by capture. Thus, many X -rays are produced, some of which radiate into space. Let n_x be the mean number of escaping X -rays produced by one incident neutrachamp, an uncertain and target dependent parameter. The neutrachamp velocity with respect to the Sun is large, $\sim 1.5 v_0$. However, the abundance of Oxygen and larger nuclei in the Sun is merely $\sim 1.4 \times 10^{-3}$. (Neutrachamp excitation off Hydrogen would require $v > 6.3 \times 10^{-3}$, well beyond what could be expected). The Moon consists of heavy elements for which the threshold velocity for neutrachamp excitation is relatively small, while the Earth's atmosphere is mostly lightweight N and O , but with a small but significant admixture of A for which $v_0 \simeq 1.0 \times 10^{-3}$. Pending appropriate estimates of n_x , and with the total champ flux given by (6.2), we deduce the secondary flux at Earth of characteristic neutrachamp X -rays from the Sun or the Moon to be $\sim 2 \times 10^{-5} n_x (\text{PeV}/M) \text{ cm}^{-2} \text{ s}^{-1}$. On the other hand, a satellite in Earth orbit should receive a much greater X -ray signal $\sim n_x (\text{PeV}/M) \text{ cm}^{-2} \text{ s}^{-1}$, uniformly distributed over the 'lower' hemisphere and due to neutrachamp encounters with the atmosphere.

Some positive champs in our Galaxy should be in the form of super-heavy atomic Hydrogen. They will display an isotope shift twice that of deuterium.

7. To Search for Ambient CHAMPS

Are champs to be found on the Earth? There may be a primordial population of terrestrial champs, but we have no way to estimate its magnitude. However, there certainly must be a flux of champs (6.2) incident upon Earth's surface if they make up the dark matter of the Galaxy. Over the 4.5 Gy age of the Earth the champ influx has amounted to $\sim 0.5 \text{ g cm}^{-2}$. Where are they to be found today? The question has both its chemical and geological connotations. Positive champ states come to rest in the atmosphere to form super-heavy Hydrogen atoms. Since the atmosphere has been Oxygen rich for $\sim 3 \text{ Gy}$, most

crust, and must be thousands of times more abundant on the Moon. Both champs and antichamps could be extracted in bulk readily and at reasonable cost. They could become a commonplace commodity. Champs could certainly find applications well beyond their cosmological utility.

They would be a delight to material science. Depending on their chemical identity, they could be used to create ultra-high density plastics or alloys. They could also render the superconducting super-collider impotent and obsolete: champs could be accelerated to modest energies and allowed to collide with an antichamp target. With a relatively small facility, physicists could study champ-antichamp annihilation at a center-of-mass energy of twice their mass, possibly exceeding the energy of any proposed or envisaged accelerator. Indeed, the annihilation process might even be used to produce energy commercially, and to justify the mining of the Moon, since the available energy content (in the form of champs and antichamps) within lunar material far exceeds that in any fissile or fusible material.

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not obligatory. Although C^+ 's alone cannot replicate the successes of the collisionless cold dark matter scenario, C^- 's (in the form of neutrachamps) do. Thus, independent searches sensitive to masses up to ~ 1000 TeV must be launched for C^+ (*i.e.*, super-heavy Hydrogen) and for C^- (higher Z 's).

An existing limit of $< 10^{-28}$ super-heavy Hydrogen atoms per proton in water [20] excludes positive champs with $M < 1$ TeV. A subsequent limit of $< 10^{-20}$ applicable to higher isotopic masses [21] could raise this lower bound to 10 TeV. However, both experiments begin with terrestrial water, which might not reflect the true abundance of super-heavy water of extreme mass. Among searches for super-heavy Carbon, [22] claims a limit of $< 2 \times 10^{-15}$ super-heavies per nucleon for $M < 100$ TeV, while [21] finds a ~ 100 times stronger limit for $M < 10$ TeV. In both cases, the specimens studied were graphite. Neutrachamps do not collapse into the galactic disk. They are not in stars, not subject to stellar processing, nor are they present in the pristine Earth. Rather, any $^{14}NC^-$ found on Earth is likely to be the result of neutrachamp exchange in the atmosphere and subsequent deposition on land and sea. The fate of this super-heavy Carbon may be to form CO_2 and to accumulate in inorganic carbonaceous minerals. On the other hand, most graphite is either crystallized magmatic material or metamorphosed sediment of organic origin. Thus, the above limits may not be relevant to the case at hand.

Similar caveats beset other searches. For example, a limit [23] of $< 1.2 \times 10^{-12}$ super-heavy Iron atoms per nucleon, valid for $M < 100$ TeV, is sensitive enough only to exclude champs with $M < 10$ TeV. A stronger limit [24] on super-heavy Lead seems irrelevant to the search for cosmic champs, since these are unlikely to capture on Bismuth.

In view of the above limits, and others, we conclude, as do Hemmick *et al.* [21], that 'it would appear safe to rule out the existence of' champs lighter than ~ 10 TeV. Curiously, the lower limit on the champ mass resulting from super-heavy isotope searches is about the same as the cosmological limits derived in Sections 4 and 5.

8. If Champs There Be

Among theorists' hypothetical effects and imagined particles, the notion of champs has a unique appeal. All of the dark matter in the cosmos is made of champs, or none of it is. It seems most unlikely that they should exist as only some of the dark matter. Thus, the champ hypothesis is readily falsified — either by searches in space or on Earth. Champs, if found, must constitute a small but non-negligible portion of the Earth's

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