Getting a charge out of dark matter

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We consider the possibility that dark matter is in the form of charged massive particles. Several constraints are discussed: (a) the absence of heavy-hydrogen-like atoms in water; (b) the agreement between the observed cosmic abundance of the elements and standard big-bang nucleosynthesis predictions; (c) the observed properties of galaxies, stars, and planets; (d) their nonobservation in γ-ray and cosmic-ray detectors, and the lack of radiation damage to space-borne electronic components. We find that integer-charged particles less massive than 10^3 TeV are probably ruled out as dark matter; but note briefly that there is a slim chance they could be blown out of the halo by supernovae. Above this mass the freeze-out abundance of these particles would overclose the Universe; thus their discovery would be evidence for inflation (or other late-time entropy dumping) below $m_{\rm ch}$. We indicate where one should consider looking for charged massive dark matter.

I. INTRODUCTION

That dark matter may be integrally charged appears at first sight entirely unreasonable. One of the fundamental properties of "dark matter" must be that it interacts only weakly with ordinary matter; the strength of the electromagnetic interaction seems to imply that dark matter is at most "millicharged." However, charged particles may comprise the dark matter if the masses are large enough that the number density required for closure is sufficiently small, and the penetration depth in matter is sufficiently large.

Discovery of stable charged massive particles (CHAMP's[1]), apart from its enormous theoretical interest, would have equally large implications for energy generation (along the lines of muon-catalyzed fusion). Another amusing, if questionable, motivation for investigating them has arisen with the recent discovery in the γ-ray burst source GB880205, of highly statistically significant absorption features at 19 and 38 keV (Ref. 6). Although explainable as cyclotron absorption features, these also correspond to the analogues of Lyman alpha lines of protons and deuterons in orbit about a massive negatively charged center. Finally, the idea of charged dark matter is a classic example of *physics under the lamp post;* for, unlike so many of the fashionable candidates, it has significant cross sections with ordinary matter and is therefore easily detected.

In this paper we consider whether the existence of ultraheavy integer-charged particles is compatible with observations and our current understanding of various astrophysical objects. Throughout this paper, unless otherwise stated, we will assume that these particles do not carry other significant interactions with ordinary matter, nor confining (i.e., strong) interactions with each other. (The stability of such massive charged particles while not automatic could be the result of some symmetry, such as, for example, if the CHAMP were the lightest supersymmetric particle in an R-parity-conserving supersymmetric theory.) Positive singly charged particles are constrained by the existing searches for heavy hydrogen up to 10^4 GeV. We propose additional constraints that raise the lower limit on the mass beyond 10^7 GeV. Discovery of such particles with masses greater than 10^7 GeV might constitute observational evidence for low-energy inflation (and hence low-energy baryogenesis), for in its absence, the hot-big-bang scenario predicts that particles this massive should have a mass density today in excess of the critical density. Some of these constraints also apply to fractional charges, and to large charges.

In Sec. II, we recall the CHAMP freeze-out abundance in a standard big-bang cosmology, indicating how the results constrain the mass, and how these constraints can be avoided. In Sec. III, we discuss the flux of CHAMP's at Earth and present the mass limit from heavy-hydrogen searches. In Sec. IV, we discuss the limits imposed by nucleosynthesis. In Sec. V we consider the effects of charged particles on galaxies, stars, and planets. In Sec.

VI we discuss limits from cosmic-ray and γ-ray data, and from the survival of space-borne electronic components. Finally, in Sec. VII we review our conclusions and indicate avenues for further investigation.

II. COSMIC CHAMP'S ABUNDANCE

Consider a charged fundamental particle $L^{\pm}$. In the early Universe, its abundance n_L is governed by the equation[7]

$$\frac{dn_L}{dt} = -3\frac{\dot{R}}{R}n_L - \langle \sigma v \rangle (n_L^2 - n_0^2) , \tag{2.1}$$

where $\langle \sigma v \rangle$ is the annihilation cross section of L^+ and L^-, R is the scale factor of the Universe, and n_0 is the number density of L's in thermal and chemical equilibrium, and we have assumed that the dominant mechanism for depletion of L is the annihilation of its own antiparticle.

We have also assumed that the number density of L^+'s and L^-'s are equal. If there is an excess of one or the other then the number density of that one will be greater than the value we compute from this equation; thus we will at least obtain a lower bound on the number density of one of the two particles. Where the presence of an excess becomes relevant we will discuss it explicitly.

For nonrelativistic L's, $\sigma_a \sim 1/v$, so $\langle \sigma v \rangle$ is a temperature-independent constant. If the L's have only electroweak interactions then, naively,

$$\sigma_a v = \frac{\pi\alpha^2}{m_L^2}N_C , \tag{2.2}$$

where N_C, the number of effective channels, is probably a number ~ 10. Following the approach of Lee and Weinberg,[7] one can calculate the energy density in L's. Expressing this as a fraction of the critical density one finds that

$$\Omega_L \approx \left[\frac{m_L}{m_0}\right]^2 \tag{2.3}$$

with $m_0 \approx 1$ TeV, for $Q_L = 1$. Since $\Omega_+ + \Omega_- \le 2$,

$$m_L \lesssim m_0 . \tag{2.4}$$

This upper bound on m_L is of scant interest, for it offers little or no actual constraint, even if the L has only electroweak interactions. For example, if L obtains its mass at the weak scale from the ordinary Higgs boson, then a 1-TeV mass implies nonperturbative Yukawa interactions, and hence a larger annihilation cross section.[8] However, we are still constrained by the limit[9] on the mass of fundamental dark-matter particles obtained by applying the Lee and Weinberg analysis to particles whose annihilation cross section is saturated by some strong interactions (i.e., $\sigma_{\text{ann}} \equiv \pi N_C / m_L^2$):

$$m_L \le m_{\max} \approx 100 \text{ TeV} h_0 N_C^{1/2} . \tag{2.5}$$

(Here h_0 is the Hubble constant in units of 100 km/s Mpc.) This upper bound on m_L may be avoided by entropy dumping, for example, an epoch of inflation at $T \lesssim T_f \sim m_L/20$, which "dilutes Ω_L" to the desired value. However, although low-energy inflationary models have been proposed,[10] they are not generic. Moreover it is unnatural to expect the inflationary expansion to have been precisely tuned such that $\Omega_L/\Omega_B \approx 0.01$.

We see that if, for charged dark matter, $m_L > m_{\max}$, we are forced into less attractive scenarios.

III. HEAVY-HYDROGEN LIMITS

The most stringent terrestrial limit on integer-charged particles comes from heavy-hydrogen searches, or more precisely from measurements of the abundance of heavy water.

The flux at Earth of L's from the galactic halo is naively

$$\phi_L \approx \frac{\rho_{\text{halo}}}{m_L} 300 \text{ km/s} \approx 9\times 10^3 \frac{\text{TeV}}{m_L} \text{cm}^{-2}\text{s}^{-1} , \tag{3.1}$$

where $\rho_{\text{halo}} \approx 0.3$ GeV cm^{-3} is the energy density of the galactic halo. We have taken the relative velocity of L and Earth to be ≈ 300 km/s, the corotation velocity at our galactic radius.

It is likely that CHAMP's would be accelerated to much higher velocities by supernova shocks, as are all other charged-particle species, and that they would remain in the galaxy after acceleration for about 10^7 yrs, as do mildly relativistic cosmic rays. While a complete review of shock acceleration, and how it generalizes to CHAMP's, is beyond the scope of this paper, we can make a rough estimate that they could be accelerated up to the same diffusion coefficient as 10^5 to 10^6 GeV protons. (Diffusivities in excess of this are difficult to attain with supernova shocks because the cosmic rays diffuse away from the shock too quickly.) Since the diffusion coefficient in a turbulent magnetic field is roughly proportional to kinetic energy per charge, it follows that CHAMP's could attain velocities as high as $c/3$ for $m_L = 10^3$ TeV.

CHAMP's may also be swept out of the solar system by the magnetic field of the solar wind. The gyromagnetic radius of a particle is

$$\rho_{\text{gyr}} = 2\times 10^{-3}\beta \left[\frac{m_L}{10 \text{ TeV}}\right]\left[\frac{\text{G}}{B}\right] \text{A.U.} \tag{3.2}$$

The magnetic field in interplanetary space in the vicinity of Earth is 5×10^{-5} G. For $\beta = 10^{-3}$, one finds $\rho \approx 0.04(m_L/10 \text{ TeV})$A.U. Thus unaccelerated, unneutralized CHAMP's (i.e., not dressed with electrons or nuclei adding up to a total charge of zero), of mass less than 250 TeV may have difficulty penetrating the solar wind. Accelerated CHAMP's, on the other hand, should have little trouble reaching Earth. Negatively charged CHAMP's have a neutral component from recombination with protons in the big bang (see Sec. IV); these will also have no difficulty penetrating the solar wind. Moreover, in a single traversal of the galactic disk, a positively charged CHAMP [or $(L\alpha)$] should attain charge-exchange equilibrium with hydrogen, and at least 1–10% of them should be neutral. At about Earth's orbit, a penetrating neutralized L^+ stands a significant chance of being photoionized by the solar UV, and can appear at Earth with a net charge. We note that neutralization of CHAMP's depends on their having integral charge. An unaccelerated CHAMP with nonintegral charge and $m/Q \lesssim 250$ TeV, may have a flux at Earth many orders of magnitude below what it is beyond the heliopause. This reduction in flux may be avoided if the CHAMP's are shock accelerated.

We next investigate whether the CHAMP's stop in the atmosphere and are incorporated into ocean water.

Consider an L^- which arrives at Earth. As we show in Sec. IV, many of these particles recombine with nuclei in the early Universe. Those that recombined with $Z \geq 2$ nuclei (mostly ^{4}He), will behave just as the L^+'s. Most of those that arrive bare, as well as the neutraCHAMP's, will pick up (or substitute) a heavy nucleus such as ^{14}N. We will discuss these further in Sec. V, where we consider constraints from atmospheric γ rays.

$L^{\pm}$'s [as well as all the (L^- nucleus) bound states, other than the neutraCHAMP's] passing through the atmosphere lose energy by Coulomb scattering off atomic electrons and off nuclei. Lindhard and Scharff have calculated the energy loss for slow massive ions in a medium. They find[11]

$$\frac{1}{\rho}\frac{dE}{dx}=\frac{1}{Am_p m_e}\left[\frac{\pi^2}{e}\frac{ZQ}{(Z^{2/3}+Q^{2/3})^{1/2}}+\frac{8\pi v}{\alpha}\frac{ZQ}{(Z^{2/3}+Q^{2/3})^{3/2}}\right], \quad (3.3)$$

where the first contribution is from atomic recoils and the second from electronic losses. For ^{14}N, the dominant component of the atmosphere, this is 370 and 80 MeV cm^2/g, respectively. Thus, the stopping distance of an L is $10^{-3}(m_L/\text{TeV})\text{g/cm}^2$; CHAMP's less massive than 10^6 TeV stop in the atmosphere's 10^3 g/cm^2 and drift down onto Earth's surface like dust.[12] The binding energy of the nuclei to the L^- is $E_B \approx 25Z^2 A$ keV, whereas the energy transfer to the bound nucleus in a scattering with atmospheric particles is at most $\beta^2 m_N \sim A$ keV. Hence the bound nuclei are not stripped from the L^-'s in the atmosphere.

The cross section for neutraCHAMP's to collide with atoms is at least $\pi/m^2\alpha^2$, leading to a $dE/\rho\, dx = 0.03$ MeV g/cm^2, although this may be increased by polarizability of the neutraCHAMP. A more reasonable estimate may be obtained by replacing m_e by m_p in (3.3), giving an atmospheric dE/dx of 0.25 MeV g/cm^2, and a stopping distance of $20(m_L/\text{TeV})\text{g/cm}^2$. Thus they too would stop in the atmosphere or ocean.

(L^+e^-) and ($L^-\alpha e^-$) are chemically identical to heavy hydrogen and, once in the ocean, are rapidly incorporated into a water molecule. Solving the diffusion equation for LHO in the approximation of an ocean of uniform temperature and density[13] we find that the drift velocity of LHO is

$$v_{\text{dr}} \lesssim \frac{m_L}{m_{H_2O}}\frac{g}{n_{H_2O}\sigma v_T} \lesssim 2\times 10^{-4}\frac{m_L}{\text{TeV}}\text{km/yr} , \quad (3.4)$$

where v_T is the thermal velocity of H_2O molecules. Given that the characteristic mixing time for the ocean is about 10 years, for $m_L \lesssim 10^3$ TeV, currents will keep the ocean well mixed. Recalling (3.1), and assuming a perfectly mixed ocean of 10 km uniform depth, we find that the number density of LHO compared to that of H_2O is

$$\frac{n_L}{n_{H_2O}} \approx 8\times 10^{-18}\frac{\text{TeV}}{m_L}\frac{t_{\text{acc}}}{\text{yr}} . \quad (3.5)$$

Here t_{acc} is the time period over which LHO accumulates in the ocean and is not removed (by chemical or other processes).

A group led by Smith[14] has obtained limits on the abundance of heavy water:

$$\frac{n_{\text{heavy}}}{n_{H_2O}} < 10^{-28}\text{--}10^{-29} \quad (3.6)$$

over the mass range 8–1.2 $10^3 m_p$, and a group at the University of Rochester (as discussed in Ref. 14) has extended these limits up to $10^4 m_p$ with a sensitivity of 10^{-24}. Moreover, they have tested not only Smith's enriched sample (about which one may perhaps entertain some doubts whether the LHO was proportionally enriched), but also ordinary sea water, up to $10^4 m_p$ with an accuracy of 10^{-16}. This contrasts with our predicted value (for $m_L = 10^4 m_p$) of

$$\frac{n_{LHO}}{n_{H_2O}} \approx 8\times 10^{-19}\frac{t_{\text{acc}}}{\text{yr}} . \quad (3.7)$$

For the LHO not to have been detected, it would have had to have been removed from the ocean with $t_{\text{acc}} \lesssim 125$ yr (40's if we use the limit of 10^{-24}). We are not aware of any mechanism for so efficiently cleansing the ocean of LHO.

Terrestrial limits thus indicate that $m_L > 10$ TeV for singly charged L^+ and ($L^-\alpha$). ($L^-\alpha$) is formed both primordially and through interactions of L^- and (L^-p), with interstellar helium. For $|Q_L| \neq 1$, the bounds are less simple to obtain. Since L^+ no longer behaves like a proton, heavy-water searches are not relevant. Similarly ($L^-\alpha$) is no longer hydrogenic; the hydrogenic complexes either contain higher Q nuclei, such as Li, or several p's and α's. Limits from the former have the difficulty that Li is less than one part in 10^9 of the baryons, so that few of the L's may be bound to them. Limits from the latter depend on the fraction of complexes containing the appropriate number of α's to be hydrogenic. Thus for $Q_L = 2$, the hydrogenic complexes are ($L^{--}ppp$) and $L^{--}\alpha p$). Given the strength of the heavy water limits, it is likely that enough of these would form so that even for $|Q_L| > 1$, we must have $m_L > 10$ TeV.

It is possible that L, in addition to being charged, also carries the quantum numbers of some confining interaction such as color. If there are stable integer-charged baryons or mesons (or even "nuclei," i.e., bound groupings of baryons or mesons), then these limits apply to those baryons or mesons.

IV. CONSTRAINTS FROM NUCLEOSYNTHESIS

The binding energy of a point nucleus of charge Z, atomic number A, to a heavy $(m_L > Am_p)L^-$ is $E_B(Z,A) \approx 25Z^2A$ keV, though, since the size of the nucleus is comparable to the radius of the innermost orbital for $Z \geq 2$, charge smearing will reduce this somewhat. In electromagnetic recombination processes (e.g., $p+e^- \rightarrow H$), the recombination temperature is[15] $T_{\rm rec} \sim E_B/40 \lesssim 0.6$ keVZ^2A (Ref. 6).

The time scale for radiative recombination of L's with nuclei of atomic number A is

$$t_{\rm rec}(A) \approx (n_A \sigma_r v)^{-1}$$
$$\approx \left[2.8\times 10^{-8} \Omega_B h_0^2 x_A 0.24 T^3 \frac{32\pi\alpha^3 Z^2 Q^4}{3\sqrt{3} m_A^2 n\beta} \right]^{-1}$$
$$\approx \left[5.1\times 10^{-5} \Omega_B h_0^2 x_A \frac{Z^2}{nA^{3/2}} \left[\frac{T}{\rm keV} \right]^{5/2} \right]^{-1} \text{ s} , \tag{4.1}$$

where x_A is the number density of A relative to baryons, n is the principal quantum number of the level to which recombination is taking place (in general[15] $n=2$), and we have taken the cross section at the recombination edge. Comparing this to the Hubble time during the radiation-dominated era, $t_H \approx (2.3\sqrt{N_F}T^2/M_{\rm Pl})^{-1}$, where N_F is the effective number of massless degrees of freedom ($N_F \approx 3.6$),

$$\frac{t_{\rm rec}}{t_H} \sim 5.8\times 10^{-3} \frac{A^{3/2} n}{x_A Z^2 Q^4} \frac{\sqrt{N_F}}{\Omega_B h_0^2} \left[\frac{\rm keV}{T} \right]^{1/2} . \tag{4.2}$$

Taking $T \approx E_B/40 = 0.6AZ^2Q^2$ keV,

$$\frac{t_{\rm rec}}{t_H} \sim 1.5\times 10^{-2} \frac{A}{x_A Z^3 Q^5} \frac{\sqrt{N_F}}{\Omega_B h_0^2} . \tag{4.3}$$

If we accept the standard big-bang nucleosynthesis results, then $0.01 \leq \Omega_B h_0^2 \leq 0.03$ and

$$\frac{0.5N_F^{1/2}A}{x_A Z^3 Q^5} \leq \frac{t_{\rm rec}}{t_H} \leq \frac{1.5N_F^{1/2}A}{x_A Z^3 Q^5} . \tag{4.4}$$

For H, ^{4}He, and ^{7}Li this comes to, respectively (with $Q=1$), $t_{\rm rec}/t_H \approx 1.0$–$2.9, 7.6$–$23, \geq 2\times 10^7$. The fraction of remaining CHAMP's which recombine with a given isotope below each recombination threshold is approximately $[1-\exp(-2t_H/t_{\rm rec})]$ implying no recombination with lithium (or deuterium, tritium, helium-3, . . .), 8–23 % with ^{4}He, and 45–68 % with H and 10–50 % remaining recombined (the numbers representing $\Omega_B h_0^2 = 0.03$ and 0.01, respectively). In fact this may be an overestimate of the fraction bound to ^{4}He since the finite-size correction to the $(L\alpha)$ binding energy means a reduced phase space for the recombination. Nevertheless, substantial fractions of the negative CHAMP's are to be found in each of three states: bare, L^-p^+), and $(L^-\alpha^{++})$.

The uniformity of ^{4}He observations in a wide range of objects, including extragalactic H II regions, leaves one confident that a bound of $n_L < n_{^4{\rm He}}$ is conservative, and even one of $n_L < 0.1 n_{^4{\rm He}}$ is reliable. Since $n_{^4{\rm He}} \approx n_B/16$,

$$m_L > 16 \frac{\Omega_L}{\Omega_B} \text{GeV} \gtrsim 0.5 \Omega_L h_0^2 \text{ TeV} . \tag{4.5}$$

Once the L's have recombined radiatively they can also undergo an exchange reaction with a higher Z or A nucleide (e.g., $L^-p+\alpha \rightarrow L^-\alpha + p$). Taking $\sigma((LA)+A' \rightarrow (LA')+A) = \mathcal{F} a_0(A)^2$, for $A' > A$, where $a_0(A) = (m_A \alpha Z_A Q)^{-1}$, we find the time scale for removing an (LA) to be

$$t_{\rm disp}^{(LA)} \approx \left[x_{A'} \Omega_B h_0^2 2.8\times 10^{-8} 0.24 T^3 \pi a_0(A)^2 \left[\frac{T}{m_A} \right]^{1/2} \right]^{-1} . \tag{4.6}$$

Comparing this to t_H,

$$\left[\frac{t_{\rm disp}^{(LA)}}{t_H} \right]_{T=E_B(A)/40} = 8\times 10^{-7} \frac{A}{QZ} \frac{N_F^{1/2}}{x_{A'} \mathcal{F} \Omega_B h_0^2} . \tag{4.7}$$

For $\mathcal{F} \sim (1)$, as expected from the adiabatic picture, most H's get displaced by α's, but few α's or ^{7}Li's get displaced; if $\mathcal{F} \lesssim 10^{-4}$ then the exchange reaction plays no role in the early Universe.[17]

Similarly, the time scale for removing an A' by the exchange reaction with an (LA) evaluated at $T = E_B(A)/40$ is

$$\frac{t_{\rm rem}(A')}{t_H} \approx 8\times 10^{-4} N_F^{1/2} \frac{(AA')^{1/2}}{QZ\Omega_L h_0^2 \mathcal{F}} \left[\frac{[L]}{[LA]} \right] \frac{m_L}{\rm TeV} . \tag{4.8}$$

Again if $\mathcal{F}$ is small the no bound is obtained, but if it is ~ 1, then for small m_L, ^{7}Li, and ^{7}Be will be depleted. Observations of [^{7}Li]/[H] in Population I stars shows a uniform upper envelope of $\sim 2\times 10^{-9}$, independent of metallicity, spectral type, etc., over a wide range of these parameters. Observations of [^{7}Li]/[H] in Population II dwarfs shows a very narrow uniform value of 10^{-10} over a significant range of spectral types, and is most plausibly attributed to primordial nucleosynthesis. Leaving aside the possibility of two unconventional mechanisms we conclude that we cannot accommodate more than 90% depletion of ^{7}Li (Ref. 18), thus

$$\frac{m_L}{\rm TeV} \geq 14 \mathcal{F} \Omega_L h_0^2 Q . \tag{4.9}$$

This, such as the ^{4}He bound, is unlikely to be more constraining than the terrestrial bound, but is also free of subtleties regarding the solar wind and galaxy dynamics.

V. CHARGED PARTICLES IN GALAXIES, STARS, AND PLANETS

If we wish to suggest that CHAMP's form the cosmological dark matter, i.e., $\Omega_{\rm ch}=1$, for which there is no convincing observational evidence, it is desirable that they also be candidates for galactic dark matter, for which there does exist evidence. We must therefore show that the time scale on which CHAMP's in the halo lose their kinetic energy and fall into the disk is long compared to the Hubble time. An important energy-loss mechanism for the CHAMP's is Coulomb scattering off electrons in the hot-ionized component of the disk interstellar medium (ISM). Energy loss for a CHAMP in this medium is $dE/dx \approx 4\pi n_e(\alpha^2/m_e v^2)\ln B$, where the Coulomb logarithm has a value of about 20. Thus $dE/dt \approx 3\times 10^{-10}(n_e\ {\rm cm}^3)(10^{-3}c/v)$ MeV/s. Since the density of the warm-ionized component of the ISM is $n_e \geq 0.5\ {\rm cm}^{-3}$ and it occupies approximately 10% of the disk volume, and since the "average" halo particle (20-kpc orbit) spends about 2% of its time in the disk, the average $dE/dt \gtrsim 3\times 10^{-13}(10^{-3}c/v)$ MeV/s. Requiring that the infall time be greater than 10^{10} yr, implies that

$$m_L \geq 10^5 \left[\frac{10^{-3}c}{v}\right]^3 {\rm TeV}\ .$$

Even allowing for uncertainties, it seems difficult to keep bare CHAMP's less massive than $\sim 10^5$ TeV in the halo.

We have neglected in the above discussion to consider the effect of shock acceleration by supernovae. It is possible that as frequently as once every 10^7 yr on average, a CHAMP is shock accelerated[20] and blown either back into the halo, or right out of the galaxy. Since the gravitational binding energy of the galaxy is approximately 10^{60} ergs, while the total energy output of galactic supernovae (assuming liberally 3×10^{51} ergs, per event, and a supernova early 10 ys) is only 3×10^{60} ergs, CHAMP ejection is marginal but not out of the question. If the reinjection into the halo is what occurs then we need only require that the infall time be greater than 10^7 yrs; the mass limit then falls to $m_L \gtrsim 100$ TeV. However, even then, we must worry about the heating of the disk matter by the CHAMP's, since this is where all that supernova energy is being deposited. The most promising scenario would seem to be that charged CHAMP's are not included in galaxies in the first place, and that they contain only neutraCHAMPs. However, as pointed out by DGS, fluctuations in which the baryons and CHAMP's are decoupled are damped too quickly to contribute to galaxy formation unless $m_L > 10^8$ TeV, and it is hard to argue that CHAMP's would be entirely excluded from galaxies. Moreover, if indeed $\mathcal{F}\sim 1$ then neutraCHAMP's would have efficiently coverted to $(L\alpha)$'s during the lifetime of the galaxy.

If despite our arguments, CHAMP's really are not contained in the galactic halo, either because they were ejected, or because they were not included in the first place, and $\mathcal{F}\lesssim 1$ so that neutraCHAMP conversion to charged CHAMP is inefficient in the galaxy, then we must reevaluate our earlier bounds. The terrestrial limit $m_L > 10$ TeV would still stand—especially if one accepts the limit of one part in 10^{24}—based on the flux of cosmological (i.e., extragalactic) CHAMP's, though that is 10^5 smaller. The limits from stars and planets derived below would be much weakened. If $\mathcal{F}\lesssim 1$ then neutraCHAMP's are probably not included in protostars, and so would not dominate the masses of stars (nor would the extragalactic CHAMP's even if they were included). CHAMP models would have effectively two types of dark matter: the neutraCHAMP's which act as cold dark matter, decoupling early from the radiation, and the charged CHAMP's, which are more strongly coupled. The ratio of their relative abundances may range from 10:1 to 1:10, depending on the relative recombination efficiencies of H and He in the early Universe.

The next issue which we consider is the effect of charged particles on stars. If the L's included in stars at the level at which they are generally present in the Universe, then they might form a significant fraction of the mass of a star, certainly distorting the very successful predictions of standard stellar evolution theory. Even at a lower abundance, they might catalyze nuclear reactions, contributing significantly to the energy production in stars.

Suppose that $L^{\pm}$'s constitute some fraction ϵ by number of particles in a protostellar cloud, i.e., $n_L=\epsilon n_B$. Then, since L's annihilate,

$$\dot{\epsilon} = -n_B\langle\sigma v\rangle\epsilon^2\ . \tag{5.1}$$

Substituting the unitarity saturated free annihilation cross section, we find

$$\tau=(n_B\langle\sigma v\rangle)^{-1}$$
$$\approx \frac{6\times 10^{23}\ {\rm cm}^{-3}}{n_B}\left[\frac{m_L}{\rm TeV}\right]^2 \frac{1}{N_C} 3\times 10^{-10}\ {\rm yr}\ . \tag{5.2}$$

The solution to (5.1) is

$$\epsilon(t)=\frac{\epsilon_0}{1+\epsilon_0 t/\tau}\ . \tag{5.3}$$

Hence, just so that the L's do not dominate the mass of the star, either we must have $\epsilon_0 \ll {\rm GeV}/m_L$, or we must have $\tau/t \ll {\rm GeV}/m_L$, the latter implying

$$\left[\frac{m_L}{\rm TeV}\right]^3 \ll 3\times 10^6\ {\rm yr}\frac{\rho_B}{{\rm g\,cm}^{-3}}\frac{t}{\rm yr}\ . \tag{5.4}$$

For $t\sim 10^9$ yr, and $\rho_*\sim 1\ {\rm g\,cm}^{-3}$ this gives naively $m_L \ll 1.5\times 10^5$ TeV. Moreover, although central concentration can enhance the annihilation rate, since almost all L^-'s inside most stars are bound in $(L\alpha)$'s the annihilation cross section is very strongly suppressed by Coulomb screening, so that the naive bound is probably conservative, and one must argue that the CHAMP's are not efficiently included in the protostar. Moreover, unless the asymmetry in L^+ vs L^- is sufficiently small, one will not be able to annihilate the CHAMP abundance down to acceptable levels anyway. Given the differing behavior of neutraCHAMP's and charged CHAMP's, such asymmetries are not unexpected.

L's which fall on stars, do not contribute to the stellar luminosity by annihilating inside them. If we wait sufficiently long, then a star will come to an equilibrium state in which all the L^+'s and L^-'s which are captured by it annihilate, so that

$$\mathcal{L}_{eq}=\pi R_*^2 \phi m_L \ . \tag{5.5}$$

This is the maximum luminosity that the L's can contribute. Comparing this to the luminosity of a blackbody at the surface temperature of the star, T_*, we find that

$$\frac{\mathcal{L}_{eq}}{\mathcal{L}_*} \approx \frac{\phi m_L \pi R_\odot}{\mathcal{L}_\odot} \left[\frac{T_\odot}{T_*} \right]^4$$
$$= 6\times 10^{-8} \Omega_L h_0^2 \frac{F}{10^5} \left[\frac{T_\odot}{T_*} \right]^4 , \tag{5.6}$$

where we have made use of Stefan's law. This is independent of $\langle \sigma v \rangle$, so that we need not worry whether or not it is the free annihilation cross section that we should be using. For stars, this ratio is indeed $\ll 1$ since stellar temperatures do not fall below $0.1 T_\odot$; however, one might feel some concern about the outer planets, which have surface temperatures below 100°. For Jupiter, Saturn, Uranus, Neptune, and Pluto this ratio is 0.6, 2, 12, 26, and 21, respectively, a fact which Fukugita, Hut, and Spergel attempted to use to explain the heating of the giant planets.[4] Hence, either the time scale for reaching this equilibrium condition must be much longer than the age of the planet, or the planet must capture only a small fraction of the incident L's (this is the case if the particles are supermassive, almost Planck mass). Assuming that the L's are perfectly mixed in the planet (i.e., they do not settle), then

$$\tau_{eq} = \sqrt{4R_*/3\phi\langle \sigma v \rangle} \ . \tag{5.7}$$

Evaluating this using the free annihilation cross section we find

$$\tau_{eq} = 7.5\times 10^{10} \text{ yr} \left[\frac{m_L}{10 \text{ TeV}} \right]^{3/2} \left[\frac{R_*}{R_\odot} \right]^{1/2} \frac{1}{\sqrt{N_C}} \ . \tag{5.8}$$

For the five outer planets this is, respectively, 24, 22, 14, 14, and 5×10^9 yr $(m_L/10 \text{ TeV})^{3/2}(1/\sqrt{N_c})$. Taking this at face value, we would conclude that for Neptune we must have

$$m_L \gtrsim \left[\frac{t_{SS}}{5.4\times 10^8 \text{ yr}} \right]^{2/3} N_C^{1/3} \ 10 \text{ TeV} \ ; \tag{5.9}$$

however, once again central concentration and Coulomb screening play competing roles to strengthen or weaken this bound.

The screening mechanism also plays an important role in the Sun, where (although as shown above, the total energy which could be emitted by annihilations is not significant) the annihilation neutrinos would have been observed unless $M_L > 10^5$ TeV (Ref. 22). L's falling onto stars will combine with ^{4}He long before they find an antiparticle against which to annihilate. Once they have done so, the Coulomb barrier will suppress the annihilation dramatically, and eliminate the neutrino signal. Recombination with ^{4}He would also suppress CHAMP catalyzation of nuclear reactions in stellar cores.

VI. COSMIC-RAY, γ-RAY, AND SATELLITE CONSTRAINTS

Although they were not designed specifically to detect CHAMP's cosmic-ray, and γ-ray detectors offer arguably the best possibility both of detecting CHAMP's in the future, and of constraining their flux levels at present.

A. Charged CHAMP's

In a cosmic-ray detector, an incident particle enters the apparatus, interacts with the material, deposits energy, and produces a signal. Some attention must be given to the form of energy deposition to which the detector responds. The exact nature of the response and hence the signal depends on the type of detector. In plastic track detectors,[23] the energy deposition results in the breaking of molecular bonds in the plastic; the resulting light molecules are etched away once the exposure is complete. In solid state detectors, energy deposition results in the formation of electron-hole pairs, which are amplified and detected electronically.

As discussed above, the energy loss of a CHAMP in a material object can be calculated from Eq. (3.3). This gives dE/dx of 500–100 MeV cm^2/g for most detectors, with the largest fraction of this coming from atomic recoils, and the rest from electronic excitations. This translates into a penetration depth of approximately $(10^3\beta)^2(m_L/10^3 \text{ TeV})$g/cm^2. The atomic recoils will result in the breaking of molecular bonds, and should cause etchable tracks in space-borne plastic detectors.[24] The loss rate for ordinary cosmic rays with $\beta > 10^{-2}$ is about $2(Z/\beta)^2$ MeV cm^2/g, and plastic track detectors respond well to $Z/\beta \gtrsim 14$. CHAMP's mimic $Z/\beta \gtrsim 16$ in loss rate, though their tracks would be peculiar in their length and profile. CHAMP's with $\beta^2 m_L \gtrsim 10^{-2}$ TeV will have penetration depths of several g cm^{-2}, and should also be detectable in balloon experiments. The approximately 80 MeV cm^2/g going into electronic excitation is sufficient to register in solid state detectors in which thresholds are approximately 3 MeV cm^2/g (Ref. 25). The likelihood of detection is much less in all detectors for neutraCHAMP's, which have substantially lower dE/dx.

One particular cosmic-ray detector of interest is the University of Chicago instrument aboard the Pioneer 11 spacecraft. The apparatus contains four separate sensors, of which we will consider only the main telescope. This instrument consists of seven elements, D1–D7, of which D1–D4 and D6 are Li-drifted silicon detectors, and D5 is a CsI scintillator viewed by a Li-drifted silicon photodiode. D7 is a plastic scintillator crystal which fits as a sleeve around the other detectors and is viewed by a photo-multiplier tube.[26] From this detector one can obtain the following limits:[25] (a) for particles which stop within 0.2 mm of Si, and lose at least 3 MeV, the flux is less than $0.01/\mathrm{cm}^2\,\mathrm{s\,sr}$; (b) for particles which penetrate at least 0.75 mm, but stop within 2.25 mm, and deposit more than 3 MeV, the flux is $\lesssim 3\times10^{-4}/\mathrm{cm}^2\,\mathrm{s\,sr}$. (c) similar limits ($10^{-2}-10^{-4}/\mathrm{cm}^2\,\mathrm{s\,sr}$) apply for all particles which stop within the telescope (a column density of 6 $\mathrm{g/cm}^2$). (d) For throughgoing particles, the flux is less than $0.3/\mathrm{cm}^2\,\mathrm{s\,sr}$.

As discussed above, the penetration depth of a CHAMP is $(10^3\beta)^2(m_L/10^3\ \mathrm{TeV})\mathrm{g/cm}^2$. For Si, with a density of $2.3\ \mathrm{g/cm}^3$ this implies penetration depths of $\gtrsim 0.4(m_L/10^3\ \mathrm{TeV})\mathrm{cm}$, for $\beta \gtrsim 10^{-3}$. Thus, for $m_L \lesssim 6\times10^3$ TeV, the CHAMP stops in the telescope and we are limited to a flux of at most $10^{-2}/\mathrm{cm}^2\,\mathrm{s\,sr}$, while for more massive (hence throughgoing) CHAMP's it can be as large as $0.3/\mathrm{cm}^2\,\mathrm{s\,sr}$. CHAMP's with $m_L \geq 6$ TeV have more than 3 MeV of kinetic energy, and lose that much in approximately 0.03 mm of Si, well less than the sizes of the individual detector elements, hence will definitely register (this is true even if one counts only the 15% of the energy coming from electronic excitations).

The naive flux of CHAMP's at Earth (i.e., neglecting any suppression due to solar wind, etc.) is approximately $1.2\times10^6/\mathrm{cm}^2\,\mathrm{s\,sr}\beta(\mathrm{TeV}/m_L)(\rho/10^{-24}\ \mathrm{g\,cm}^{-3})$, giving us a limit of

$$\frac{m_L}{10^3\ \mathrm{TeV}} > 1.2\times10^2(10^3\beta)\frac{\rho}{10^{-24}\ \mathrm{g\,cm}^{-3}} \tag{6.1}$$

for $m_L < 6\times10^3$ TeV. This implies that the density at Earth of CHAMP's with $m_L < 6\times10^3$ TeV is $\rho < 8\times10^{-27}\ \mathrm{g/cm}^3(10^{-3}/\beta)(m_L/10^3\ \mathrm{TeV})$. This is less than the local dark-matter density of $10^{-24}\ \mathrm{g/cm}^3$. Suppression of the flux by the magnetic field in the solar wind, which is limited for L^+ by resonant charge exchange in the interstellar medium, could perhaps reduce the flux by 1 or 2 orders of magnitude for $m_L \lesssim 10^2$ TeV, but this is insufficient to relax these limits, CHAMP's are to be the halo dark matter. These bounds do not apply to neutraCHAMP's, which are not detected in cosmic-ray detectors.

For $m_L > 6\times10^3$ TeV, the limit is

$$\frac{m_L}{10^3\ \mathrm{TeV}} > 4(10^3\beta)\frac{\rho}{10^{-24}\,\mathrm{g\,cm}^{-3}} \tag{6.2}$$

which is not constraining. The flux limits on CHAMP's can be improved by using the fact that the background from ordinary cosmic rays can be resolved into charge peaks. Better limits might also be obtained using pulse-height analysis, since, unlike the usual cosmic rays, the energy loss for CHAMP's does not fall strongly as they lose energy.

The overall flux limits from plastic track detectors are quite severe, with the exposure being measured in $m^2\mathrm{sr\,yr} = 3\times10^{11}\ \mathrm{cm}^2\mathrm{sr\,s}$. This would be sufficient to place a limit of $m_L > 3.5\times10^{17}\beta(\rho/10^{-24}\ \mathrm{g\,cm}^{-3})$ TeV, pushing the CHAMP mass up to the grand-unified-theory (GUT) scale. Even if one is skeptical of using the entire exposure to limit the CHAMP flux, one can still require that the CHAMP flux be smaller than that of the cosmic rays (which are resolved into charge peaks with $\delta Z/Z \ll 1$). Taking, for example, the CHAMP flux to be less than that of fluorine ($3\times10^{-4}\ \mathrm{cm}^2\mathrm{s}$), one finds that

$$m_L \gg 3\times10^7\ \mathrm{TeV}\frac{\rho}{10^{-24}\ \mathrm{g\,cm}^{-3}}\ . \tag{6.3}$$

Once again, the bound may be weakened because it does not apply to neutraCHAMP's, but only to charged CHAMP's. Moreover, plastic track detectors, have not yet been calibrated at $\beta \approx 10^{-3}$; one may require $\beta \gtrsim 10^{-2}$ to have a distinguishable signal. If one uses the whole exposure, this mild amount of acceleration seems unlikely to suppress the flux by the 12 orders of magnitude necessary to allow CHAMP's less massive than the $m_{\max}$; but it might provide the 5 orders of magnitude necessary if one accepts only the weaker bound (6.3). Despite these reservations, results from plastic track detectors are definitely problematic for CHAMP's lighter than $m_{\max}$.

One type of cosmic-ray detector, which we do not generally classify as such, is the electronic components of satellites. Since CHAMP's cause molecular bonds to break, they will degrade electronic components. The energy deposited by CHAMP's per mass of an exposed object is

$$\mathcal{E} \approx 3.5\times10^3\beta\frac{\mathrm{TeV}}{m}\frac{\rho}{10^{-24}\mathrm{g\,cm}^{-3}}\left[\frac{Z}{7}\right]^{2/3}\ \mathrm{Mrad/yr}\ ,$$

whereas the observed damage is[27] less than 2 krad/yr. This gives $m_L \gtrsim 3\times10^3(\rho/10^{-24}\ \mathrm{g/cm}^3)$TeV. Even a factor of 10 reduction in flux due to the solar wind magnetic field still leaves a useful bound.

In conclusion, given the bounds from the solid-state and plastic track detectors, and from the satellites, it us unlikely that charged particles less massive than $m_{\max}$ comprise the galactic halo.

B. NeutraCHAMP's

When a negatively charged CHAMP, or a neutraCHAMP strikes Earth's atmosphere, it has a significant probability of combining with a heavy nucleus. For the CHAMP this is a radiative recombination, while for the neutraCHAMP it is an exchange process. In Sec. IV, we presented the cross sections for both these processes. In the atmosphere the mean free path for a neutraCHAMP to replace its proton with a N^{14}, is $(\rho\lambda)_{\mathrm{ex}} \approx 1.6\times10^6\beta$ $\mathrm{g/cm}^2$ while for a CHAMP to pick up an ^{14}N, $(\rho\lambda)_{\mathrm{rec}} \approx 4\times10^{10}\beta^2\ \mathrm{g/cm}^2$. Here we have used the Born approximation (as suggested by DGS) to calculate the exchange cross section. Since we expect this to be an underestimate of the cross section (by several orders of magnitude) the bound we obtain will be conservative. However, if the charge-exchange cross section is as large as the geometric cross section of the neutraCHAMP, then most neutraCHAMP's will have converted to $(L\alpha)$'s during the lifetime of the galaxy, and the charged CHAMP limits apply.

For both the radiative and the charge exchange processes, we expect ~ 1 of the recombinations to be excited states of $(L^-\,{}^{14}\mathrm{N})$, and to be followed by the emission of a cascade of few hundred keV to 3.5 MeV γ rays. The atmospheric attenuation length of such photons is (10–30) $\mathrm{g/cm}^2$, so we expect $(0.6-2)\times10^{-5}/\beta$ of the photons from the neutraCHAMP's, and $(2.5\text{–}7.5)\times 10^{-10}/\beta^2$ of the CHAMP photons to be observable at any given point in the atmosphere if the emission occurs uniformly in the atmosphere. For the neutraCHAMP's this is probably a good approximation, at least for interesting masses $m_L \gtrsim 10^3$ TeV, because of the low-energy-loss rate. The charge CHAMP's stop in $(m_L/10^3\ \text{TeV})(10^3\beta)^2\ \text{g/cm}^2$, and so the photons will be emitted predominantly in the upper atmosphere. Considering the neutraCHAMP's we expect a γ-ray flux of

$$\frac{d\phi_\gamma}{d\Omega} \approx 8\frac{\text{TeV}}{m_L}\frac{\rho L_p}{10^{-24}\ \text{g cm}^{-2}} g_\gamma(E)/\text{cm}^2\text{s sr}\ , \qquad (6.4)$$

where $g_\gamma(E)$ is the photon multiplicity. The observed limits on the flux of MeV γ rays are[28]

$$\frac{d\phi_\gamma}{d\Omega} \lesssim 2\times 10^{-4}/\text{cm}^2\text{s sr} \qquad (6.5)$$

in narrow line features, and

$$\frac{d^2\phi_\gamma}{d\Omega\, dE} \lesssim 5\times 10^{-5}/\text{cm}^2\text{s keV sr} \qquad (6.6)$$

in the continuum in this energy range. The lines (2–4-keV resolution detectors) thus give

$$\frac{m_L}{\text{TeV}} \gtrsim 6\times 10^4 \frac{\rho L_p}{10^{-24}\ \text{g cm}^{-3}} g_\gamma(E) \qquad (6.7)$$

while in the continuum with 100 keV resolution detectors we get

$$\frac{m_L}{\text{TeV}} \gtrsim 2.5\times 10^3 \frac{\rho L_p}{10^{-24}\ \text{g cm}^{-3}} g_\gamma(E)\ . \qquad (6.8)$$

Thus even in the event that there happened to be no lines in the previously investigated wavebands (especially near 511 keV), one would still find that $m_L >$ few 10^3 TeV. Bounds on charged CHAMP's are about an order of magnitude less, due to the longer recombination mean free path.

In fact, one anticipates even stronger bounds coming from detailed analysis of the dependence of the γ-ray background on height in the atmosphere, and on zenith angle.

Cosmic- and γ-ray detector data constitute the most serious constraints to the whole idea of charge dark matter. Because they put stringent limits on both the charged CHAMP's and the neutraCHAMP's, one is forced into scenarios where the CHAMP abundance is not its thermal freeze-out abundance, but has been diluted by late time entropy dumping.

VII. CONCLUSION

We have seen that the possibility that the particle making up the dark matter is integerly charged, while intriguing, has definite difficulties. If CHAMP's are to be the cosmological dark matter, we would like them to also be good candidates for galactic dark matter. Moreover, we would prefer that they be lighter than $m_{\max} \lesssim 10^3$ TeV, the maximum mass of a fundamental dark-matter particle whose abundance is determined by the usual freeze-out mechanism, and does not require fine-tuned dilution. This possibility has essentially been ruled out.

In Fig. 1 we summarize the various bounds which we have discussed.

For a singly charged particle, the terrestrial limits on heavy water tell us that the mass of the particle must be greater than 10 TeV. Weaker, though somewhat more general bounds are obtained by consideration of the abundances of the light elements, in particular ^{4}He. Consideration of the infall of CHAMP's from galactic halos into the disk implies that $m_L \gtrsim 10^5$ TeV for charged CHAMP's, if they are to make up the galactic halo. It is possible, that charged champs are ejected from the disk and either reinjected into the halo, or blown right out of the galaxy, by shock acceleration in supernovae. The en-

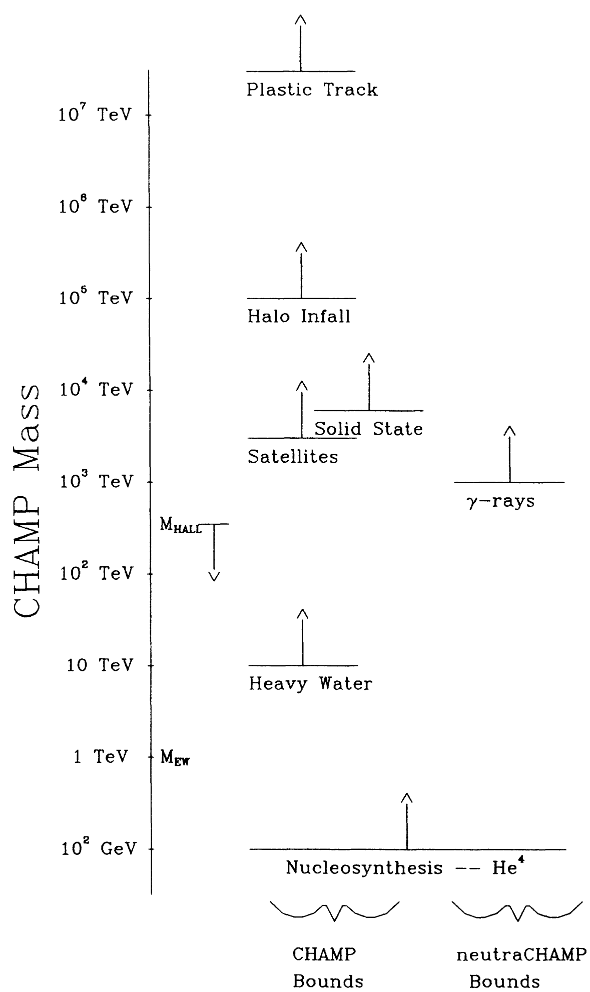

FIG. 1. Limits on the CHAMP mass obtained in this paper. The bounds have been separated into those applying only to charged CHAMP's, and those applying to neutraCHAMP's $m_{\max}$ has also been displayed, as has M_{ew}, the mass at which a particle with only electroweak couplings would naively have the appropriate freeze-out density to provide the dark matter.

ergetics of the former possibility are marginal, while in the latter possibility one must still be concerned about the heating of the disk. It is unlikely that galaxies do not contain charged CHAMP's *ab initio*, since, unless $M_L \gtrsim 10^8$ TeV, decoupled CHAMP-baryon fluctuations damp on excessively short time scales.

Cosmic-ray and γ-ray detectors further diminish the attractiveness of CHAMP dark matter. Data from the Pioneer 11 silicon telescope imply that charged CHAMP's lighter than 6×10^3 TeV cannot make up the halo dark matter. Consideration of data from plastic track detectors place a limit on charged CHAMP's of $m_L \gtrsim 3\times10^{14}$ TeV $(10^3\beta)(\rho/10^{-24}\ \mathrm{g\,cm^{-2}})$, if one uses the entire exposure, and $m_L > 3\times10^7(\rho/10^{-24}\ \mathrm{g\,cm^{-2}})$ TeV, if one uses only the flux of fluorine cosmic rays as the limiting flux. Even with the closure density of 10^{-29} $\mathrm{g\,cm^{-3}}$, this would force $m_L \gtrsim m_{\max}$, unless there is a substantial suppression of the charged CHAMP flux near Earth.

Data from γ-ray observations strengthen the case that m_L must be larger than $m_{\max}$. They show that for neutraCHAMP's $m_L \gtrsim \mathrm{few}\times10^3$ TeV, if they make up the halo dark matter. Of course, it is possible that CHAMP's constitute only the cosmological dark matter and that the galactic dark matter is something completely different. This is not a particularly appealing possibility.

Since non-integer-charged particles cannot be neutralized by recombination with ordinary matter, the charged CHAMP bounds appropriately adjusted for the different charge apply to them, as does the possibility that they have been ejected from the galaxy and evade the bounds.

We see that if we are to consider CHAMP dark matter seriously, we are forced to consider masses greater than $m_{\max}$, thus giving up our hope that its abundance will be determined by simple freeze-out considerations. Although this is not entirely fatal, it is discouraging. Nevertheless, in some sense it gives us even greater motivation to search for CHAMP's, since detection of them would be evidence for inflation, not to mention providing a catalyst for low-temperature fusion.

As we have discussed, cosmic-ray and γ-ray detectors are both well suited in searching for charged dark matter. They are also complementary since one looks for charged CHAMP's best and the other for neutraCHAMP's. Other possibilities include heavy isotope searches in moon rocks, antarctic ice, pieces of asteroids, or any other materia. Looking for reduced mass shifts in atomic lines would also be of interest. Given the comparative simplicity of the task compared to the search for more weakly interacting dark matter, and given the possibility of simultaneously detecting the dark matter and obtaining evidence of inflation, a search for charged ultramassive particles seems well worth the effort.

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[1]While this work was in preparation, it came to our attention that charged dark matter is under independent consideration by De Rújula, Glashow, and Sarid (Ref. 2) from whom we adopt the acronym CHAMP for charged massive particles, and the name neutraCHAMP for neutral bound states of CHAMP's and nuclei. Ideas arising from conversations with them or from preliminary versions of their work will be acknowledged in each case. CHAMP's have also been discussed previously though in different contexts by Glashow and Cahn (Ref. 3) and by Fukugita, Hut, and Spergel (Ref. 4). Some of the issues in this paper have also been addressed independently by Chivukula and Walker (Ref. 5) in the context of technicolor cosmology.

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[16]Peebles (Ref. 15) performed this calculation for the specific case of protons and electrons recombining to form hydrogen atoms. The case of positively charged nuclei recombining with L^-'s is fundamentally the same. Necessary to Peeble's result were the justification of several assumptions, mostly involving the legitimacy of taking the photon spectrum and the particle kinetic energy distributions to be thermal over relevant energy ranges during the recombination process. We have verified that all the necessary approximations are also valid for the case of nuclei and L^-'s. Taking equations (26) and (30) of Peebles, we can then see that recombination occurs at $T_r \simeq E_B \ln[n_N(2\pi/m_N T_r)^{3/2}] \simeq E_B/40$, when photoionization becomes suppressed relative to recombination.

[17]De Rújula, Glashow, and Sarid (DGS) attempted to calculate the displacement cross section explicitly for $LP+\alpha \rightarrow L\alpha + p$; using the Born approximation they find $\sigma_{\mathrm{exch}} \approx (1/v)2^{19/4}\pi^{3/2}\alpha^{17/4}(m_p r_0)^{9/4}Z'^{5/4}[G(Q)/m_p^{1/2}Q^{3/2}]$ with $r_0=1.2$ fm, $G(Q)=z/[\exp(z)-1], z=2\pi(Z'-1)\alpha(m/2Q)^{1/2}$, and $Q=E_B(A')-E_B(A)$. Evaluating this at $T=E_B(A)/40$, when the channel opens, this gives $\mathcal{F}((LA)A' \rightarrow (LA')A) \approx 6\ 10^{-4}G(z)(A/A')^{3/2}ZZ'^{-7/4}$, so that $(t_{\mathrm{disp}}^{(LA)}/t_H) \gg 1$. However, the validity of the Born approximation is questionable since the velocity of $A' \ll \alpha$, so that the wave function of the incoming A' is significantly altered in the interaction region. Thus, for example, in elastic electron-hydrogen scattering, the Born approximation fails entirely for electron energies less than 100 eV.

[18]We should caution that we have not considered the effects of the L on the nucleosynthesis itself. Since the binding energies of L^- with the various nucleides are as large as various nuclear binding energies, they can act as catalysts and may drastically affect important reaction rates, and hence the outcome of standard big-bang nucleosynthesis (SBBN). Thus there may exist a region of particle and cosmological parameter space in which the observed element abundances are well reproduced, but which is not related to the region of parameter space in which ordinary SBBN work. Such possibilities have been discussed previously, for example, in the context of late decaying particles (Ref. 19). The problem here however is more difficult as the nuclear cross sections involve are unmeasurable (unless we manage to find some L^-'s), and such as most nuclear reaction rates, probably uncalculable.

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