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DARK MATTER: COSMOLOGICAL LIMITS AND APPLICATIONS

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ABSTRACT

The astrophysical consequences and limits from particle theory on dark matter are reviewed. A discussion of some observational consequences is included.

The nature of the primary constituent of the Universe is of obvious interest to both cosmologists and particle theorists. With regard to cosmology, it appears that the bulk of the Universe must be in some form of non-luminous matter, i.e. dark matter (DM). Whatever the DM may be it poses several problems[1)] on various cosmological distance scales. If inflation is correct, then we know in addition that at least some of the DM is something other than baryons. Thus we demand from the particle theory a candidate. At the same time however, for almost every candidate, there are several cosmological constraints on its mass lifetime, couplings etc. In this contribution, I will review the DM problems and discuss the cosmological constraints on the DM candidates. I will also discuss some of the implications of having a galactic halo filled with heavy ($m >$ few GeV) DM.

One of the chief problems concerning DM is that as one looks at increasing distance scales there appears to be more and more non-luminous matter. The amount of DM is generally described in terms of the overall cosmological density in DM

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$$\Omega = \rho/\rho_c = \rho/(1.88\times10^{-29}h_o^2 \text{ g cm}^{-3}) \tag{1}$$

where ρ is the average mass density of DM and ρ_c is the critical mass density to close the Universe and $h_o = H_o/100$ km Mpc^{-1} s^{-1} is the Hubble parameter. On a given scale, Ω is determined via a mass-to-light ratio,[2)]

$$\Omega = (M/L)\mathcal{L}/\rho_c \tag{2}$$

where $\mathcal{L} = 2\times10^8 h_o\ L_\odot/Mpc^3$ and $L_\odot$ is a solar luminosity. In the solar neighborhood, $(M/L) \simeq 2$ and $\Omega \sim 10^{-3}/h_o$. Mass to light ratios for the inner luminous parts of spiral galaxies yield $\Omega \simeq 10^{-2}$. On the scale of binaries and small groups of galaxies (which would include galactic halos) $\Omega \simeq 0.05 - 0.15$. On the largest scales on which determinations of Ω exist[3)], Ω is no larger than a few tenths. Finally, inflation requires $\Omega = 1$. Hence, the hierarchy of DM problems.

What this DM hierarchy implies is that as we go to larger scales, there appears to be more DM, i.e. the M/L ratios continue to increase. Already on galactic scales, there is good evidence from rotation curves[2)] of spiral galaxies for the presence of dark matter and a galactic halo. The rotation curve is a measure of the velocity as a function of distance from the center of the galaxy of a star as it revolves around the galaxy. If there were no DM, one would expect that at distances beyond the bulk of the luminous matter that $v^2 \sim 1/r$. Instead one finds flat rotation curves ($v^2 \sim$ constant) out to very large distances ($\gtrsim 50$ kpc). This implies that the mass of the galaxy must continue to increase $M \sim r$ beyond the luminous region.

As a first guess as to the identity of the DM, one might pick baryons i.e., ordinary matter. From big bang nucleosynthesis, there are good limits[4)] on the value of Ω in the form of baryons. In the standard model, one finds good agreement for the predicted abundances of the light isotopes D, ^{3}He, ^{4}He, ^{7}Li only for a range in Ω_B

$$0.01 \lesssim \Omega_B \lesssim 0.15 \tag{3}$$

For $\Omega_B < 0.01$, D and ^{3}He are overproduced while for $\Omega_B \gtrsim 0.15$, ^{4}He is overproduced and D is underproduced. If inflation is correct and $\Omega = 1$, then at least some of the DM must be non-baryonic.

On a galactic scale, baryonic DM is consistent with nucleosynthesis. There are however numerous arguments[5] against a baryonic halo. Put briefly, it is very difficult to have a large baryon density in such a way that it is unobservable. In the form of gas the baryons would heat up and emit X-rays in violation of observed limits. To put the baryons in non-nuclear burning stars (Jupiters) would require an extrapolation of the stellar mass distribution which is very different from what is observed. Dust or rocks along with dead remnants such as neutron stars or black holes would require a metal abundance in great excess of the galactic metallicity. Very massive ($\gtrsim$ 100 $M_\odot$) black holes remain a possibility.

There are of course, many other candidates for the DM. Because of its important role in the formation of galaxies, DM has classified[6] into three types: hot, warm and cold DM. They are distinguished by their effective temperature at the time they decoupled from the thermal background. Examples of hot particles are neutrinos or very light Higgsinos with $\lesssim$ 100eV masses. These particles decouple at $T_d \sim 1$ MeV and are thus still relativistic at T_d. Warm particles decouple earlier and have higher masses (up to ~ 1 keV). Any superweakly interacting neutral particle is a warm candidate such as a right handed neutrino. Cold particles are non-relativistic at temperatures relevant for galaxy formation and usually have masses $\gtrsim$ 1 GeV. Examples of these include heavy neutrinos, photinos/Higgsinos, sneutrinos and axions. With this classification, the specific identity of the particle is no longer important. The benefits and problems associated with each type of DM with regards to galaxy formation has been nicely reviewed in the contribution of M. Davis[7] and the reader is referred there for further details.

Given the need for DM, we can ask what sort of constraints are there. The most common cosmological constraint is on the mass of a stable particle and is derived from the overall mass density of the Universe. The mass density of a particle X can be expressed as

$$\rho_x = m_x Y_x n_\gamma \lesssim \rho_c \simeq 10^{-5} h_o^2 \ \mathrm{GeV/cm^3} \tag{4}$$

where $Y_x = n_x/n_\gamma$ is the density of x's relative to the density of photons $n_\gamma = 400(T_\gamma/2.7°K)^4$ for $\Omega h^2 < 1$. Hot particles have limits characteristic to that of neutrinos. For neutrinos[8] $Y_\nu = 3/11$ and one finds

$$\sum \left(\frac{g_\nu}{2}\right) m_\nu \lesssim 100 \ \mathrm{eV} \ (\Omega h^2) \tag{5}$$

where the sum runs over neutrino flavors and $g_\nu = 2$ for a Majorana mass neutrino and $g_\nu = 4$ for a Dirac mass neutrino. It has also been pointed out recently[9] that if two of the neutrino flavors are unstable, their decays will presumably produce at least one of the lighter neutrinos so that the cosmological mass limit on it would be $m_{\nu_L} \lesssim 30\mathrm{eV}(\Omega h^2)$. All hot particles with abundances Y similar to neutrinos will have mass limits as in eq. (5).

Warm particle limits are derived from eq. (4) as well. Warm particles have lower abundances than neutrinos and the corresponding mass limits are weaker. Recall that $Y_\nu = 3/11$ is derived from the conservation of entropy before and after $e^\pm$ annihilation. Neutrinos at this time are decoupled so that after the annihilations $(T_\nu/T_\gamma)^3 = 4/11$ and $Y_\nu = (3/4)(T_\nu/T_\gamma)^3 = 3/11$. (The factor of 3/4 is due to the difference between Fermi and Bose statistics.) If a particle x interacts more weakly than neutrinos then the ratio $(T_x/T_\gamma)^3$ will be lowered[10] due to other particle species' annihilations. Thus Y_ν is reduced allowing[11] for a larger value for m_x. If the particle x decouples around the GUT epoch, then Y_x could be as low as $O(10^{-2})$ and $m_x \lesssim O(1)$ keV.

For cold particles the analysis is somewhat different. The abundance is now a function of m_x and in most cases one finds a lower limit to m_x. The reason for this is that for large m_x, Y_x is controlled by the annihilations of x. When the annihilations freezeout, Y_x is fixed. The freezeout will then depend on the annihilation cross-section and roughly one finds $Y_x \sim (m_x\sigma_A)^{-1}$ and $\rho_x \sim 1/\sigma_A$. This situation was first analyzed for neutrinos[12]. The annihilation cross-section in this case is basically $\sigma_A \sim m_\nu^2/m_W^4$ so that $\rho_X \sim 1/m_\nu^2$ and yields[12,13] $m_\nu \geq 2$ GeV for Dirac mass neutrinos and $m_\nu \geq 6$ GeV for Majorana mass neutrinos.[14,13]

Supersymmetric theories introduce several DM candidates. The reason is that if the R-parity (which distinguishes between "normal matter and the supersymmetric partners) is unbroken then there is at least one supersymmetric particle which must be stable. Candidates for the stable particle include the photino, Higgsino, sneutrino, gravitino, and goldstino. If we assume for simplicity that all of the scalar quarks and leptons have equal masses then the photino annihilation cross-section can be expressed as[15,16,17]

$$\langle\sigma v\rangle_A = \frac{8\pi\alpha^2}{m_{sf}^4} \sum_f q_f^4 \, (1-z_f^2)^{1/2} \, m_{\tilde\gamma}^2 \, (z_f^2 + 2x(1-17z_f^2/8)) \tag{6}$$

where α is the fine structure constant, m_{sf} is scalar fermion mass, q_f the electric charge of the fermion f, $z_f = m_f/m_{\tilde\gamma}$ and $x = T/m_{\tilde\gamma}$. For $m_{sf} \simeq 40$ GeV, $m_{\tilde\gamma} \geq 1.8$ GeV. For Higgsinos[17], the annihilations are controlled by the fermion Yukawa couplings and the cosmological bound requires $m_{\tilde H} \geq m_b$ or about 5 GeV.

Sneutrinos are an interesting example in that there is in general no cosmological limit on their mass.[18] In addition to the standard weak annihilations of sneutrinos, there is also the process $\tilde\nu + \tilde\nu \to \nu + \nu$ via zino exchange. In this case $\langle\sigma v\rangle_A \sim 1/M_{\tilde Z}^2$ and is independent of $m_{\tilde\nu}$. Thus $\rho_{\tilde\nu}$ is fixed by parameters other than $m_{\tilde\nu}$ making the sneutrino mass free from cosmological bounds. Before turning to the gravitino, I note that the goldstino arguments are essentially warm particle limits, and will depend on its specific interactions.

The remaining possible supersymmetric DM candidate is the gravitino. Although in most models the gravitino is not stable, there is nothing which prevents its stability. If stable, the gravitino mass limit would again be $m_{3/2} \leq O(1)$keV as for a warm particle[19] assuming that its abundance $Y_{3/2}$ was determined by considering gravitino decoupling at the Planck time. Such an early decoupling will make $Y_{3/2}$

sensitive to the details of inflation.[20] In general, gravitinos will be produced in the reheating after inflation so that[17,21]

$$Y_{3/2} \simeq (2\text{-}3) \times 10^{-3}\, T_R/M_P \tag{7}$$

where T_R is the temperature to which the Universe reheats and $M_P = 1.22 \times 10^{19}$ GeV is the Planck mass. Thus for a standard value $m_{3/2} \simeq 100$ GeV one must have $T_R \lesssim 10^{12}$ GeV to satisfy the cosmological bound (eq. 4).

If the gravitino is unstable, there are various interesting cosmological constraints which can be applied depending on what the gravitino decays into. The possible decay channels include, $\gamma + \tilde{\gamma}$, $\nu + \tilde{\nu}$ or goldstone boson and goldstino. For simplicity, I will assume that there is only one supersymmetric particle lighter than the gravitino. The decay rate of the gravitino into photinos has been computed[21]

$$\Gamma = m_{3/2}^3/4M_P^2 \tag{8}$$

for $m_{3/2}$ well above threshold. The decays into scalar and fermion should not be that much different and I will use eq. (8) for all decays.

Gravitino decays into $\nu,\tilde{\nu}$ pairs appears to be the least interesting cosmologically, as it would be very difficult to observe the neutrino. The neutrino energy today is given in terms of the temperature at which the gravitino decays T_D and $m_{3/2}$.

$$E_\nu = \frac{m_{3/2}}{2}\left(\frac{T_\gamma}{T_D}\right) \tag{9}$$

where $T_\gamma = 2.7°K$ is the present temperature of the microwave background. The ratio (T_D/T_γ) is determined by setting the lifetime of the gravitino $\tau_{3/2} = \Gamma^{-1}$ equal to the age of the Universe

$$t = t_o(T_\gamma/T)^{3/2} \tag{10}$$

where $t_o = 2\times10^{17}h_o^{-1}$ s. Thus

$$(T_D/T_\gamma) \approx 60h_o^{-2/3}m_{3/2}^2(\mathrm{GeV}) \tag{11}$$

and

$$E_\nu \approx 8\times10^{-3}h_o^{-2/3}/m_{3/2}(\mathrm{GeV}) \quad \mathrm{GeV} \tag{12}$$

Hence for $m_{3/2} = 1-100$ GeV, $E_\nu \approx 100$ keV - 10 MeV. Clearly for this range for neutrino energies it would be impossible to pick out the decay neutrinos from the solar background. Only if the sneutrino and gravitino were somewhat degenerate so that the rate eq. (8) is reduced could one boost E_ν.

Depending on the sneutrino mass there is always the limit on T_R due to the mass density of sneutrinos coming from the decay,

$$\rho_{\tilde{\nu}} = m_{\tilde{\nu}}\, Y_{3/2}\, n_\gamma \tag{13}$$

or

$$T_R \lesssim 10^{14}/m_{\tilde{\nu}}(\mathrm{GeV}) \quad \mathrm{GeV} \tag{14}$$

and is not a particularly strong limit.

Gravitino decays into a goldstone-goldstino pair such as an axion and axino give no cosmological limit on T_R but can provide a scenario for galaxy formation.[22] Among the problems of neutrino dominated models (see ref. 7), is that the scale on which clustering occurs is too big. One way to shorten this scale is to consider a decaying particle scenario[23] in which the Universe becomes matter dominated earlier (because of the increased mass of the particle relative to a 30eV neutrino for example) and density perturbations begin to grow earlier on a smaller scale. The decay has in addition the advantage of distributing a large fraction of the mass of the Universe uniformly (something which is not possible in standard cold scenarios but is possible in "biased" scenarios, see ref. (7)). All of these models however require a late decay, i.e., one in which $T_D/T_\gamma \lesssim 5\text{-}10$ implying that the decay rate $\Gamma \sim 10^{-40}$ GeV. Such a small rate is characteristic of gravitational decays $\Gamma \sim G_N\, m^3 \sim 6\times10^{-39} m^3$ GeV making[22] the decay of a gravitino into a goldstone-goldstino pair a plausible candidate for this type of scenario.

The final channel for the gravitino is into a photon, photino pair. This possibility has been the subject of a lot of attention by many groups and I will here only summarize the results as given by ref. 24). The constraints all place limits on the abundance $Y_{3/2}$ or the reheating temperature T_R and come from a wide variety of sources ranging from big bang nucleosynthesis to the cosmic microwave background. The weakest limit is due to the mass density during nucleosynthesis and requires $T_R \lesssim 10^{18}/m_{3/2}$ GeV ($m_{3/2}$ is in GeV). Entropy production after nucleosynthesis requires $T_R \lesssim 2\times10^{13}\, m_{3/2}^{1/2}$ GeV. Similar to the limit in eq. (14) there is a limit due to the mass density of photinos produced in the decay $T_R < 10^{14}/m_{\tilde\gamma}$ GeV. Requiring that the decay of the gravitino does not over deplete the ^{4}He abundance through photodissociation means that $T_R \lesssim 5\times10^{14}/m_{3/2}$ GeV. From the isotrotry of the microwave background radiation one finds $T_R \lesssim 5\times10^9\, m_{1/2}^{1/2}$ GeV and finally the most stringent limit is due to the photoproduction of D + ^{3}He and implies that $T_R \lesssim 2.5\times10^{10}/m_{3/2}$ GeV. In terms of a limit on $Y_{3/2}$, we have from this last constraint $Y_{3/2} \lesssim 2.5\times10^{-12}/m_{3/2}$.

If the supersymmetric particle spectrum is such that $m_{\tilde\gamma} < m_{3/2}$ the above constraints greatly restrict inflationary models. It is interesting to note however that simple models of inflation in the context of N=1 supergravity naturally satisfy this bound. I am referring to models[25-27] in which inflation is described by a superpotential and a single chiral superfield ϕ such that $f(\phi) = \mu^2 g(\phi)$ and all couplings in $g(\phi)$ are O(1). Thus there is only one scale associated with inflation namely, μ. The parameter μ can be determined from the magnitude of density perturbations produced during inflation

$$\frac{\delta\rho}{\rho} \sim 10^3\, \mu^2 \sim 10^{-4} - 10^{-5} \qquad (15)$$

The reheating temperature is also determined by μ

$$T_R \sim \mu^3 M_p \tag{16}$$

Thus a value for $\mu \sim 10^{-4}$ from eq. (15) implies $T_R \sim 10^7$GeV and $Y_{3/2} \sim 10^{-15}$. Although such a low value of T_R tends to make baryon generation more difficult[29] by requiring low values for the masses of Higgs fields which violate baryon number conservation, the mass of the ϕ field is $m_\phi \sim \mu^2 M_p \sim 10^{11}$GeV and this is just about the bound[30] on the Higgs field mass, allowing for a sizeable baryon asymmetry.

If we assume that $Y_{3/2} \sim 10^{-15}$ and the decay rate in eq. (8) for gravitinos to $\gamma + \tilde{\gamma}$, then the decay is capable of producing a feature[31] in the γ-ray background with energy equal to the neutrino energy in eq. (12). Hence for a gravitino mass $m_{3/2} \sim O(10)$ GeV, the γ-ray energy will be $\sim O(1)$ MeV with a flux comparable to the observed bump at ~ 1 MeV.

In the remaining part of this contribution, I will discuss some of the consequences of having a galactic halo consisting of cold DM. As was discussed earlier, there are several arguments against having baryonic DM in the galactic halo.[5] I will assume therefore for the remainder of this paper that the DM consists of nonbaryonic matter. Although there is no real reason against hot or warm DM in galactic halos, what I am about to discuss only applies to cold DM.

The eventual verification of the existence of a cold DM galactic halo obviously depends on some kind of signature or signal. One interesting suggestion[32] has been to use a very cold detector with superconducting grains which would flip as the DM passes through. What I will discuss here is some possible signatures due to DM annihilations in the halo and in the sun. I will throughout assume that there is one type of DM with $\Omega \simeq 1$, and a local density $n_x = (0.3/m_x)\ \mathrm{cm}^{-3}$ with velocities $v \simeq 300\ \mathrm{km\ s}^{-1}$.

The possible observation of annihilation of DM in the galactic halo was recently examined[33] in the case where photinos of mass $m_{\tilde{\gamma}} \simeq 3$ GeV were the DM. The annihilations of photinos can lead to appreciable fluxes of cosmic rays. Although the γ-ray flux from these annihilations is well below the established backgrounds the predicted flux of positrons and antiprotons is significant. The predicted fluxes are $F_{e^+} \approx 5\times10^{-4}$ cm^{-2} s^{-1}sr^{-1} and $F_{\bar{p}} \approx 1.5\times10^{-6}$ cm^{-2}s^{-1}sr^{-1} while the observed fluxes are[34] $F_{e^+} \approx 10^{-3}$cm^{-2}s^{-1}sr^{-1} and $F_{\bar{p}} \approx 3\times10^{-6}$cm^{-2}s^{-1}sr^{-1}. This is particularly important for the case of the antiproton flux because these are low energy $\bar{p}$'s ($0.6 \lesssim E \lesssim 1.2$ GeV) whose origins are otherwise difficult to explain.

The effects of cold DM in the sun were first discussed by Steigman etal.[35] and more recently in the context of the solar neutrino problem.[36-38] Annihilations in the sun[39] can however lead to an appreciable flux of high energy neutrinos. The flux of neutrinos depends primarily on the rate of capture by the sun of the DM which I assume to be a photino here for definiteness. The capture rate is computable[37,39] in terms of the photino-proton elastic scattering cross-section σ_E

$$\Gamma_c \approx 10^{29}\sigma_{E,36}/m_{\tilde{\gamma}} \ s^{-1} \tag{17}$$

where $\sigma_{E,36} = \sigma_E/10^{-36}\text{cm}^2$. In order to have $\Omega = 1$ the annihilation cross-section must be $\langle\sigma v\rangle_A \approx 10^{-26}\text{cm}^3\text{s}^{-1}$ so that the annihilation rate in the sun is[39]

$$\begin{aligned}\Gamma_A &\approx \frac{4\pi}{3} R_\odot^3 n_{\tilde{\gamma}}^2 \langle\sigma v\rangle_A \\ &\approx 2\times10^{54}(n_{\tilde{\gamma}}/n_p)^2 \ s^{-1}\end{aligned} \tag{18}$$

where $n_{\tilde{\gamma}}/n_p$ is the ratio of the number density of photinos to protons. This yields an equilibrium number density

$$n_{\tilde{\gamma}}/n_p \approx 2\times10^{-13} \ \sigma_{E,36}^{1/2}/m_{\tilde{\gamma}}^{1/2} \tag{19}$$

Photinos in the sun with the abundance given by eq. (19) would also elastically scatter inside the sun providing a means for transporting energy and possibly resolving the solar neutrino problem.[35-37] The rate for energy transport is

$$\dot{E}_T = n_{\tilde{\gamma}} \langle\sigma v\rangle_E \, \Delta E N_p \tag{20}$$

where $\Delta E \sim m_{\tilde{\gamma}} v^2 \sim$ few keV and $N_p \sim 10^{57}$ is the number of protons in the sun. Energy transport is significant when $\dot{E}_T \sim \mathcal{L}_\odot = 4\times10^{33}$ erg s^{-1} or

$$\langle\sigma v\rangle_E (n_{\tilde{\gamma}}/n_p) \sim 10^{-38} \text{ cm}^3 \text{ s}^{-1} \tag{21}$$

Using eq. (19) this requires

$$\sigma_{E,36} \sim 0(100) m_{\tilde{\gamma}}^{1/3} \tag{22}$$

What one finds however is that typically[39,38] $\sigma_{E,36} \sim 0(10^{-2})$ and eq. (21) cannot be satisfied by any known DM candidate such as heavy neutrinos, photinos, Higgsinos or sneutrinos. The primary reason being the large annihilation rate resulting in the low abundance (eq. 19).

It is just these annihilations however that may lead to a signature for the DM in the galactic halo.[39] Whatever the DM is, annihilations will lead to high energy neutrinos whose flux is given by

$$\phi_\nu = \frac{1}{2} N_\nu \Gamma_C / 4\pi (1\text{A.U.})^2 \tag{23}$$

$$= 16 \, N_\nu \sigma_{E,36} / m_x \text{ cm}^{-2}\text{s}^{-1}$$

where N_ν is the number of neutrinos produced per annihilation. The flux in eq. (23) must then be compared to the background flux of neutrinos which is produced by cosmic ray collisions in the atmosphere. If we consider the production of electron and muon neutrinos, the solar and atmospheric fluxes are shown[39] in Table 1.

Table 1

X	$\nu_{e\odot}$	ν_e ATM	$\nu_{\mu\odot}$	ν_μ ATM
Dirac mass neutrino; m_ν=5 GeV				
E_ν=200MeV	0.08	.3-1.6	0.16	0.7-3.8
E_ν=5GeV	3×10^{-3}	9×10^{-5}	3×10^{-3}	5×10^{-4}
Majorana mass Neutrino m_ν=12GeV				
E_ν=600MeV	0.11	0.03-0.08	0.22	0.09-0.21
Photino				
$m_{\tilde{\gamma}}$=5GeV,E_ν=200MeV	0.27	0.3-1.6	0.54	0.7-3.8
$m_{\tilde{\gamma}}$=10GeV,E_ν=500MeV	0.08	0.06-0.14	0.16	0.14-0.35
Higgsino				
m_H=5GeV,E_ν=200MeV	0.5	0.3-1.6	1.0	0.7-3.8
Sneutrino				
$m_{\tilde{\nu}}$=5GeV,E_ν=5GeV	0.25	9×10^{-5}	0.25	5×10^{-4}

all fluxes in $cm^{-2}s^{-1}$

$$\nu_e = \nu_e + \bar{\nu}_e \; ; \; \nu_\mu = \nu_\mu + \bar{\nu}_\mu$$

The ranges for the atmospheric neutrino fluxes[40)] corresponds to the differences associated with latitude and up-going and down-going fluxes due to geomagnetic effects. As one can see from Table 1, in most cases the solar flux of high energy neutrinos is at least comparable to that of the atmospheric flux making their detection possible. Finally, it has been noted[41)] that for the cases of a Dirac mass neutrino or a sneutrino, the earth traps the DM and these annihilations within the earth also lead to a strong source for high energy neutrinos.

In conclusion, there is a broad region of overlap between cosmology, astrophysics and particle physics with respect to DM. The need for DM comes from several sources; inflation; galaxy formation; galactic halos, etc. Supersymmetry is thus of great interest in that it most probably guarantees one stable new particle. Indeed combining theory with new experimental results seems to require[42)] that $\Omega_{\tilde{\gamma}} h_o^2 \geq$ 0.0025 (recall the limit for baryons[4)] is only $\Omega_B h_o^2 \geq 0.01$). The consequences of a positive detection of DM in the halo are enormous. A single galaxy formation scenario may be singled out and if supersymmetric the identity of the lightest and stable SUSY partner would greatly narrow the choice of supersymmetric models.

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