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**COSMOLOGICAL AND ASTROPHYSICAL ASPECTS OF THE EARLY
UNIVERSE**

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Cosmological and Astrophysical Aspects of the Early Universe

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Various applications of particle physics to astrophysics and cosmology are reviewed, with special reference to cosmic rays as probes. Inflation is discussed in connection with the COBE data, and indications for a combination of hot and cold dark matter are mentioned. Several particle candidates for the dark matter are introduced, including neutrinos as hot dark matter and the lightest supersymmetric particle as cold dark matter. Some possible signatures are discussed, including neutrino oscillations in the light of solar neutrino experiments, and the annihilations or elastic scattering of massive particles in the galactic halo. Among the cosmic ray particles of interest are neutrinos, photons, muons, positrons, antiprotons and the lightest supersymmetric particles.

1. INFLATION ?

This is the most primordial application of particle physics to cosmology that we will discuss here, and perhaps the most far-reaching. It is also one which may now be entering the realm of experimental test with the apparent discovery [1] by COBE of fluctuations in the microwave background radiation. Inflation has been proposed [2] as a solution to many of the most fundamental problems in cosmology. One of the most acute is the homogeneity and isotropy of the Universe on distance scales much larger than the event horizon. Distant parts of the Universe that are only now coming within the present horizon $d = c \cdot t_0$ were seemingly almost identical at the epoch of matter-radiation decoupling when they were separated by $O(100)$ event horizons. Another major problem is the flatness of the Universe, corresponding to the fact that the present density of the Universe is close to the critical value that marks the boundary between an open and a closed Universe. This means that the curvature term $\propto k$ in the Hubble expansion law

$$H^2 \equiv \left(\frac{\dot{a}}{a}\right)^2 = \frac{8\pi G_N}{3}\rho - \frac{k}{R^2} \quad (1)$$

is at most comparable to the matter density term in which $\rho \propto R^{-3}$ or 4 , which appears unnatural unless the curvature k vanishes. A related problem is the great age and size of the Universe. A generic solution of the general relativistic equations for a Robertson-Walker-Friedmann Universe should have a scale factor a of order

the Planck length $l_P = 10^{-33}$ cm, and an age of order the Planck time $t_P = 10^{-43}$ s. Finally, the primary motivation for inflation was in fact the need to dispose of unwanted massive particles that might have been produced in the very early Universe, such as magnetic monopoles [3].

All these problems would be solved if there was an early epoch of exponential expansion of the Universe driven by an effective cosmological constant Λ :

$$H^2 = \left(\frac{\dot{a}}{a}\right)^2 = \frac{\Lambda}{3} \Rightarrow a(t) \simeq \exp(Ht) : \\ H = \sqrt{\frac{\Lambda}{3}} \quad (2)$$

In such a case, the true event horizon would have expanded far beyond the naive $d = c \cdot t$, and parts of the Universe that are now becoming visible would actually have been within the event horizon before the inflationary epoch as seen in Fig. 1, so that the mystery of their similarity disappears.

The first inflationary models [2],[4] postulated the existence of a scalar field ϕ stuck in a false vacuum $\langle \phi \rangle = \phi_0$ with large positive energy, giving an effective cosmological constant

$$\Lambda = 8\pi G_N V(\phi_0) \quad (3)$$

Unfortunately, in these old inflationary models the phase transition would never have been completed [5], because the parts of the Universe in between the bubbles of true vacuum would have expanded much more rapidly than the bubbles themselves. The first solution to this difficulty was to make the bubbles expand exponentially

Fig. 1 - The expansion of the Universe in conventional Big Bang cosmology and in inflationary cosmology (bold lines), showing how the latter resolves the horizon problem [2].

themselves, by postulating a slow rollover towards the true vacuum [6] as in Fig. 2b. Getting a long enough rollover to accommodate the present Universe and have sufficiently small density primordial density perturbations requires a very flat potential, which is most naturally realized in the context of supersymmetry [7]. Subsequently, other solutions to the difficulties of old inflation have emerged, such as chaotic [8] and extended [9] inflationary models, but their details do not concern us here. For most inflationary potentials, gravitational waves would be much less important in the COBE data than density perturbations [10], but this is not a universal feature.

The density perturbations in new inflationary models are related to quantum fluctuations that result in different speed of rollover in different parts of the Universe. They are usually almost scale-independent [11]:

$$\delta\rho/\rho \propto \left[\ln(m_p \cdot d) \right]^p \quad (4)$$

because of the near time-independence of the inflationary de Sitter phase.

Fig. 2 - (a) Old inflation [2] occurred while the Universe was stuck in a local minimum of the potential, to be followed by tunnelling to the true vacuum.

(b) New inflation [6] occurs during a slow roll-over into the true vacuum.

A cause of concern to us is the fact that no field-theoretical potential suitable for inflation has ever been derived from string theory [12]. Our hunch is that string theory may do the job of inflation in a different way. We find that string theory near the cosmological singularity resembles a Hall conductor [13], and speculate that the density perturbations may be related to calculable Hall model correlation functions that could perhaps even be measured in the laboratory [14].

The COBE data [1] reveal to us for the first time primordial density perturbations when they were small and in the linear growth regime, free of the uncertainties of astrophysical evolution. COBE is really breaking new ground, taking us for the first time back to an epoch in the history of the Universe prior to cosmological nucleosynthesis. One of the key points to watch in the future will be whether the microwave background fluctuations are Gaussian, as predicted by inflation but not by some other theories of structure formation.

2. HOW MUCH DARK MATTER ?

Inflationary models all predict that the density of the observable Universe is very close to the critical density: $\Omega \simeq \Omega_c \equiv 1$. However, observations

give much less visible matter: $\Omega_{vis} \simeq 0.01$, whilst Big Bang nucleosynthesis and the LEP measurement [15] of N_ν also indicate a low baryon density: $\Omega_{bar} < 0.1$. Galactic rotation curves indicate that there is more matter in the galactic halo: $\Omega_{halo} = 0.1$, which could mathematically be baryonic, e.g. “Jupiters”, but not gas, dust or “snowballs”. However, there are arguments [16] against these possibilities, and physically it seems more reasonable to expect that some of whatever dark matter makes the global density up to the critical value also appears in the galactic halo, whose local density here is estimated [17] to be

$$\rho_{halo} \simeq 0.3 \text{ GeV cm}^{-3} \quad (5)$$

with an uncertainty of perhaps a factor of 1.5.

The question then arises whether this dark matter is hot, i.e. relativistic at the epoch of decoupling, or cold, i.e. non-relativistic at that epoch. If structures in the Universe were seeded by a Gaussian random field of perturbations, as in models of inflation, then, with hot dark matter alone, structures would not have formed early enough on small distance scales (see Fig. 3a), whereas cold dark matter would have favoured equally the growth of fluctuations on all distance scales (see Fig. 3b) [18]. On the other hand, small seeds such as loops of cosmic string would have given too much power at short distances if combined with cold dark matter, and would be less unsatisfactory when combined with hot dark matter [19].

New light on the nature of the dark matter has been cast by the COBE data. By themselves they do not favour hot or cold dark matter, but when combined with previous data at smaller comoving distance scales, such as those on cluster-cluster correlations, a model which contains mainly cold dark matter but with an admixture of hot dark matter to give more power at large scales seems to be preferred [20] (imagine adding Figs. 3a and 3b). The best fit seems to be when

$$\Omega_{cold} \simeq 0.7, \quad \Omega_{hot} \simeq 0.3, \quad \Omega_{bar} \lesssim 0.1 \quad (6)$$

but a pure cold dark matter model cannot yet be excluded. Whilst there are many candidates for the cold dark matter component, by far the most plausible candidate for the hot dark part of the matter is a neutrino weighing about 10 eV.

Fig. 3 - Impression of the spectrum of mass fluctuations due to scale-invariant Gaussian perturbations with (a) hot dark matter, (b) cold dark matter.

3. PARTICLE DARK MATTER CANDIDATES

3.1. Relic Neutrino Density

There are three mass regions where a neutrino or other neutral particle may have a cosmologically interesting relic density, as seen in Fig 4. When $m_\nu \ll 1$ MeV, the relic number density is essentially independent of m_ν , and hence the relic mass density is almost linear in m_ν . It increases past the critical density when the single neutrino mass is about 30 eV [21]. After reaching a peak far above ρ_{crit} , it turns down and falls below the critical density again when m_ν is a few GeV [22], essentially as a result of a Boltzmann suppression factor in the number density. Somewhere above $M_Z/2$, the density rises again [23] past the critical density, because the annihilation cross-section is falling like m_ν^{-2} so that the number density is no longer falling.

Recent data on solar neutrinos can be construed as supporting the hypothesis of a hot dark matter neutrino in the first of these cosmologically interesting regions. If the apparent neutrino deficits now seen by four experiments [24] are ascribed to particle physics rather than some modification of the standard solar model [25], the preferred fits to the data (see Fig. 5) have

$$\Delta m^2 \simeq 10^{-5} \text{ eV}^2, \quad \sin^2 \theta \simeq 10^{-2} \text{ or } 1 \quad (7)$$

The smallness of the neutrino masses is understood by the GUT seesaw mechanism [26] as being due to mixing with very massive conjugate neutrinos:

Fig. 4 - Impression of the relic neutrino density as a function of its mass, illustrating three regions where the critical density could be reached: $m_\nu \sim 30 \text{ eV}$, 3 GeV , 1 TeV .

Fig. 5 - Regions of the $(\sin^2 2\theta, \Delta m^2)$ plane in a two-flavour matter-enhanced neutrino oscillation interpretation of the available solar neutrino data [24].

$$(\nu, \nu^c) \begin{pmatrix} \simeq 0 & m_D \\ m_D & m_M \end{pmatrix} \begin{pmatrix} \nu \\ \nu^c \end{pmatrix}, \quad (8)$$

where the off-diagonal Dirac mass term m_D is often a charge-2/3 quark mass, and M_M may be as large as M_{GUT} . The seesaw mass matrix (8) favours

$$m_{\nu_e} : m_{\nu_\mu} : m_{\nu_\tau} \propto m_u^2 : m_c^2 : m_t^2, \quad (9)$$

and hence

$$m_{\nu_e} \ll m_{\nu_\mu} \simeq 10^{-2\frac{1}{2}} \text{ eV} \ll m_{\nu_\tau} \simeq 10 \text{ eV} \quad (10)$$

as also seen in Table 1, in which case the ν_τ is a natural candidate to be the hot dark matter. If one assumes a texture for the neutrino mass mixing matrix that is similar to that for quarks, one expects that

$$\sin^2 2\theta_{e\mu} \simeq 10^{-2}, \quad \sin^2 2\theta_{\mu\tau} \simeq 10^{-3} \text{ to } 10^{-4}. \quad (11)$$

The former is consistent with the solar neutrino data, and the latter would be detectable [27] with the forthcoming generation of laboratory neutrino oscillation experiments [28] (CHORUS, NOMAD, FNAL P803). Moreover, ranges of neutrino mass differences approaching the solar neutrino range ($\Delta m^2 = 10^{-4} \text{ eV}^2$, $\sin^2 2\theta > 10^{-3}$) can be probed [29] by sending a CERN neutrino beam towards the ICARUS detector in the Gran Sasso laboratory, or towards SuperKamiokande, as seen in Fig. 6. It just so happens that the SPS

Fig. 6 - Regions of the $(\sin^2 2\theta, \Delta m^2)$ plane already probed with accelerator and reactor neutrino experiments, and accessible to the ICARUS or Superkamiokande detectors using CERN neutrino beams [29].

Table 1

Values of the $\nu_\mu - \nu_\tau$ mixing parameters, τ -neutrino mass, and relic cosmological density as a function of the top-quark mass and central values of the parameters (including $M_3/M_2 = 10$). Where relevant we have assumed that $m_{\nu_\mu} = 2 \times 10^{-3}$ eV.

m_t (GeV)	$\sin^2 2\theta_{\mu\tau}$	$\Delta m_{\nu_\mu - \nu_\tau}^2$ (eV ²)	m_{ν_τ} (eV)	$\Omega_\nu h^2$
90	1.7×10^{-3}	22	4.7	0.05
100	1.3×10^{-3}	38	6.2	0.07
110	9.8×10^{-4}	66	8.1	0.09
120	7.5×10^{-4}	115	11	0.12
130	5.6×10^{-4}	202	14	0.16
140	4.2×10^{-4}	368	19	0.21
150	3.0×10^{-4}	715	27	0.29

to LHC proton beam transfer line points conveniently towards the Gran Sasso, whilst the SPS North Area beam extraction points conveniently towards Japan!

3.2. Implications for Dark Matter of LEP Measurements of N_ν

It is well known that LEP can constrain invisible Z^0 decays by measuring precisely the resonance line-shape. This is because the hadronic peak cross-section is given by

$$\sigma_{had}^{peak} = \frac{12\pi}{m_Z^2} \left(\frac{\Gamma_{ee}}{\Gamma_Z} \right) \left(\frac{\Gamma_{had}}{\Gamma_Z} \right). \quad (12)$$

where the effective number of neutrino species N_ν is defined by

$$\Gamma_Z = N_\nu \Gamma_\nu + 3\Gamma_e + \Gamma_{had}. \quad (13)$$

The peak cross-section (12) changes by 13 % for $\Delta N_\nu = 1$, so one can hope to measure N_ν with a precision better than 0.1 if one has adequate statistics and can control the collision luminosity with a precision better than 1 %. This is the case at LEP, and indeed they find [15] together that

$$N_\nu = 3.04 \pm 0.04 \quad (14)$$

Thus there are certainly only three light neutrino species, but also massive dark matter candidates are severely constrained [30],[31]. Indeed, the measurement (14) allows a massive

conventionally-coupled Dirac neutrino only if

$$m_{\nu_D} > 44 \text{ GeV} \quad (15a)$$

a Majorana neutrino only if

$$m_{\nu_M} > 38 \text{ GeV} \quad (15b)$$

and a massive sneutrino only if

$$m_{\tilde{\nu}} > 42 \text{ GeV} \quad (15c)$$

Going back to Fig. 4, we see that the relic massive neutrino density must be very small unless $m_\nu > M_Z/2$:

$$\Omega_{\nu_D} < 2 \times 10^{-5}, \quad \Omega_{\nu_M} < 2 \times 10^{-3} \quad (16)$$

which rules these dark matter candidates out [30] unless they weigh much more than $M_Z/2$.

3.3. Upper Limit on the Mass of a Neutrino-like Dark Matter Particle

Returning now to the third region of mass in Fig. 1, we find that the present mass density

$$\Omega_X \simeq \frac{m_X \cdot n_X}{2 \times 10^4 h^2 T_0^2}, \quad (17)$$

where n_X is the relic number density of the candidate X , h is the Hubble constant in units of $100 \text{ km s}^{-1} \text{ Mpc}$, and T_0 is the black body temperature. From this one can infer that

$$\Omega_X h^2 \simeq \frac{1}{10^3 \langle \sigma v \rangle \text{ TeV}^2} \quad (18)$$

where $\langle \sigma v \rangle$ is the velocity-averaged annihilation cross-section. Hence, parametrizing

$$\langle \sigma v \rangle \simeq C \cdot \frac{\alpha^2}{m_X^2} \quad (19)$$

one finds that [23]

$$m_X \simeq (16 \propto \sqrt{c}) \sqrt{\frac{\Omega_X h^2}{0.25}} \text{ TeV} \lesssim 1 \text{ TeV} \quad (20)$$

In other words, such a dark matter particle cannot weigh much more than 1 TeV, and could even be a "guaranteed" discovery for the LHC or SSC!

3.4. Lightest Supersymmetric Particle

Most supersymmetric theories contain one stable particle which could be a cosmological relic. This is because they usually have a multiplicatively conserved quantum number called R-parity [32] which takes the values +1 for all particles and -1 for all sparticles. The conservation of R-parity is linked to that of baryon number B and lepton number L , since

$$R = (-1)^{3B+L+2S} \quad (21)$$

This also shows how R-parity could be violated, by violating either B or L conservation, but we assume for now that R-parity is indeed conserved. This has three important consequences: 1) sparticles are always produced in pairs, 2) heavier sparticles decay into lighter sparticles, and 3) the lightest supersymmetric particle (LSP) is stable.

In this case, the LSP is likely to be neutral and only weakly interacting [33]. Otherwise, it would presumably have condensed along with ordinary baryonic matter and bound to it to make anomalous heavy isotopes. These have not been seen at abundance levels far above that expected for such a relic LSP [34], so we assume that it does not bind, and hence must be electromagnetically neutral and only weakly-interacting. There are candidates to be the LSP with spin 0 - the sneutrinos, spin 1/2 - some gaugino/Higgsino mixture, and spin 3/2 - the gravitino.

We have already found the bound (14c) on $m_{\tilde{\nu}}$. Moreover, *sneutrinos* would usually have had coherent interactions with heavy nuclei, which have not been seen [35], telling us that

$$m_{\tilde{\nu}} \lesssim 8 \text{ GeV or } > 1 \text{ TeV} \quad (22)$$

Furthermore, some sneutrinos could have been captured by the Sun and annihilated there to yield high-energy solar neutrinos, as discussed in more detail in the next section. From their absence we infer [36] that

$$m_{\tilde{\nu}} \lesssim 3 \text{ GeV} \quad (23)$$

In view of all these conflicting bounds, we conclude that the LSP is not a sneutrino.

Gravitinos would give "warm" dark matter if they had the canonical number density and weighed about 1 keV so as to give the critical density [37]. Such a value of the gravitino mass is not favoured theoretically, and one usually finds that the gravitino weighs more than the selectron and sneutrino, and hence is not the LSP. Let us hope so, because the interactions of the gravitino are only of gravitational strength, so that it would be essentially unobservable.

Neutralinos are spin-1/2 mixture of the four neutral gauginos and Higgsinos in the minimal supersymmetric extension of the Standard Model. They are parametrized by a bare gaugino mass $m_{1/2}$, a Higgs mixing parameter μ , and the ratio $\tan \beta$ of Higgs vacuum expectation values. The content of the LSP simplifies in the following limits:

$$m_{1/2} \rightarrow 0, \quad LSP \simeq \tilde{\gamma};$$

$$\mu \rightarrow 0, \quad LSP \simeq \tilde{H}^0, \quad (24)$$

but a generic LSP is a complicated mixture of states. Experiments at LEP and CDF constrain a neutralino LSP to weigh more than about 10 to 20 GeV [31],[38]. Even with this constraint, there are large regions of parameter space where the LSP can have the critical density, as seen in Fig. 7 [39].

On the basis of the above discussion, we conclude that the LSP is a very plausible candidate for the cold dark matter, and we will now discuss how to look for it.

4. POSSIBLE WAYS TO DETECT SUPERSYMMETRIC RELICS

4.1. Annihilation in the Galactic Halo

Annihilations of relic particles in the galactic

Fig. 7 - Upper limits on the halo density of relic particles from the non-observation of products from annihilations in the galactic halo [42].

halo would yield fluxes of stable particles in the cosmic rays given by

$$\phi_{\bar{p}, e^+, \gamma} \propto n^2 \sigma \tau \quad (25)$$

where n is the density of relic particles in the halo, σ is their annihilation cross-section, and τ is the time that the annihilation product spends in the halo. There used to be an old claim [40] to have observed low-energy antiprotons in the cosmic rays that could be due to relic annihilations. However, this has been contradicted by more recent data [41], and realistic calculations [42] of the particle fluxes, that fix the annihilation cross-section by requiring that it lead to the critical density, use the estimate (5) of the density of dark matter in the galactic halo, and the Lund model to estimate secondary particle spectra, predict antiproton fluxes that are lower than even the new upper limits. They also find that positrons do not give useful constraints on the supersymmetric model parameters, as seen in Fig. 7.

It has been suggested [43] that LSP annihilations might yield an observable monochromatic γ line. Unfortunately, it seems that this would require a fine adjustment of supersymmetric model parameters that does not seem very natural.

4.2. Annihilation inside the Sun

Cold dark matter particles in the galactic halo passing by the neighbourhood of the Sun have some probability of interacting with one of its nuclei while they pass through the Sun, losing some energy to nuclear recoil, and getting trapped [44] in an elliptical orbit with a perihelion below the solar surface. The trapping probability is given by

$$\Gamma_T \propto \pi R_\odot^2 n_{\text{relic}} \bar{v} \sigma_{\text{elastic}}, \quad (26)$$

and, after trapping, repeated collisions with solar nuclei give the dark matter particles a quasithermal distribution in the core of the Sun. Their density is controlled by evaporation, rather like the escape of light gases from planetary atmo-

Fig. 8 - Upper limits on the halo density of relic particles from the non-observation of high-energy solar neutrinos from annihilations in the Sun [48].

spheres [45], if the relic particles weigh less than about 3 GeV , which is unlikely to be the case for the LSP. The density of heavier relics is controlled by mutual annihilation at a rate

$$\Gamma_A = \frac{1}{2} \Gamma_T \quad (27)$$

Neutrinos can be produced directly by annihilations, or by second and third generation particles that yield secondary neutrino fluxes from prompt decays, which can be estimated using the Lund Monte Carlo [46]. Comparisons with upper limits from underground experiments looking for neutrinos or muons (produced by neutrino interactions in rock) from the direction of the Sun seem to rule out sneutrinos as dark matter candidates. However, they seem [47] to be consistent with the LSP being a neutralino, as seen in Fig. 8. However, searches are already starting to constrain supersymmetric model parameters, and offer better hope for detecting LSPs than does the previous search for free halo annihilations.

One important ambiguity in the interpretation of the searches for such high-energy solar neutrinos is the possibility that they might be subject

to Mikheyev-Smirnov-Wolfenstein (MSW) matter oscillations [48]. These could either enhance the ν_e flux and suppress the $\nu_\mu + \nu_\tau$ flux, as seen in Fig. 9a, or vice versa, as seen in Fig. 9b. This means that the conservative way to interpret high-energy solar neutrino searches, until their possible oscillation parameters are pinned down, would be to assume that they could have oscillated into the flavour for which your detector has the smallest detection efficiency, probably ν_τ . Eventually, if a high-energy solar neutrino flux is actually seen and its production mechanism is understood, measurements of the fluxes of different flavours could be used to pin down the oscillation parameters, and even distinguish between different explanations of the apparent low-energy solar neutrino deficit [48].

4.3. Elastic Scattering in the Laboratory [49]

Halo relic particles travel with a velocity v which is typically of the order of 300 km s^{-1} , and hence a collision with a nucleus will deposit a recoil energy which is typically

Fig. 9 - Examples how matter-enhanced neutrino oscillations might (a) increase, (b) decrease the flux of $\nu_\mu + \nu_\tau$ at the Earth due to relic annihilations in the Sun [49].

$$\Delta E < m_\chi v^2 \simeq 10 \left(\frac{m_\chi}{10 \text{ GeV}} \right) \text{ keV} \quad (28)$$

The effective interaction lagrangian is [50]

$$\mathcal{L}_{eff} = \sum_q \left[(\bar{\chi} \gamma_\mu \gamma_5 \chi) \cdot (\bar{q} \gamma^\mu (A P_L + B P_R) q) + (\bar{\chi} \chi) \cdot C m_q \bar{q} q \right], \quad (29)$$

where $P_{L,R}$ are chirality projection operators and A, B, C are model-dependent coefficients. The first term in (29) is spin-dependent and due to sfermion and Z exchange, whereas the second term is spin-independent and due to Higgs and sfermion exchange. The nucleon matrix elements of the spin-dependent interaction

$$M_{SD}(\chi_n^p \rightarrow \chi_n^p) = 4 \bar{s}_\chi \underline{s}_n^p \sum_{q \in p} (B - A) \Delta q \quad (30)$$

are controlled by the contributions Δq of the different flavours of quarks to the nucleon spin, which have been measured by experiments at CERN [51] and elsewhere [52]. The spin-independent matrix elements [53]

$$M_{SI}(\chi_n^p \rightarrow \chi_n^p) = \left[\frac{\hat{f} m_u C_u + m_d C_d}{m_u + m_d} + f C_s + \frac{2}{27} (1 - f - \hat{f}) (C_c + C_b + C_t) \right] \quad (31)$$

are controlled by the contributions of the different flavours of quarks to the proton and neutron masses.

Going from nucleonic matrix elements (30), (31) to nuclear matrix elements requires a model of the nucleus, and the simplest and first to be tried is the naive single-particle version of the shell model, according to which the angular momentum of an odd-even or even-odd nucleus is carried by the single unpaired nucleon. The accuracy of this model can be assessed by comparing its predictions for nuclear magnetic moments with data [54]. It works reasonably well for light nuclei that are near to a closed shell, but looks unreliable for many heavier nuclei of interest to dark matter searchers. A better model is the odd-group model [55], according to which the nuclear angular momentum is carried collectively by all the nucleons of the odd species. When this model can be checked, for example by using mirror pairs of nuclei or by complete shell model calculations [56], it seems to accurate to within 10% or so, which should be adequate for our purposes.

Figure 10 shows for many different nuclei a figure of merit for the spin-dependent interaction which includes kinematic and odd-group model factors [57]. Among the lighter nuclei, F seems to be favoured, also because the odd-group model is likely to be rather reliable in this case. Among the heavier nuclei, Tl has a large estimated figure of merit, though this cannot be considered as reliable as for F. Some elements often used in previous underground experiments, such as Ge and Si, do not have a very favourable spin figure of merit. However, spin-independent interactions are likely to be important for heavier nuclei, particularly if the Higgs mass is small and/or $\tan \beta$ is large [27].

The total elastic scattering rate is

Fig. 10 - Figure of merit for the spin-dependent contribution to relic scattering off different nuclei [59].

$$R = (R_{SD} + R_{SI})Y \left(\frac{4m_\chi m_N}{(m_\chi + m_N)^2} \right) \left(\frac{\rho_\chi}{0.3 \text{ GeV cm}^{-3}} \right) \left(\frac{\langle v_{rel} \rangle}{320 \text{ km s}^{-1}} \right) \text{ events kg}^{-1} \text{ d}^{-1} \quad (32)$$

where Y is the abundance of any isotope considered, and the spin-dependent contribution is

$$R_{SD} = 9.5 \left(\frac{10^4 \text{ GeV}^2}{m_f^2 - m_\chi^2} \right) \lambda^2 J(J+1) \zeta(r_{spin})$$

$$\left(\frac{m_f^2 - m_\chi^2}{e^2} \sum_q (B - A) \Delta q \right)^2 \quad (33a)$$

where λ is determined by the odd-group model and $S(\Gamma_{spin})$ a spin form factor, whilst the spin-independent contribution is

$$R_{SI} = 210 m_N^2 m_\chi^4 \zeta(r_{charge}) \left(\hat{f} \frac{m_u C_u + m_d C_d}{m_u + m_d} + f C_s + \frac{2}{27} (1 - f - \hat{f}) (C_e + C_b + C_t) \right)^2 \quad (33b)$$

where $\zeta(r_{charge})$ is a charge form factor. To estimate the scattering rate, we assume that the cosmological density of the LSP is critical, and adjust the scalar mass or $\tan\beta$ accordingly. Figure 11a shows the regions of the $(\mu, m_{1/2})$ plane where the rate is estimated to be above 0.01 events per day for elastic scattering off a 1 kg target that is half ^{73}Ge and half ^{76}Ge [57]. The rates are very similar for the two halves of the target, and a kg of NaI would give a similar rate [58], as seen in Fig. 11b. Figures 11c,d show analogous calculations for F and Tl targets. Figure 12 presents the Ge results in a different way, showing the “probability” of dark matter detection calculated as the fraction of the parameter space in which the critical density can be attained for which the elastic scattering rate is above 0.1 events per kg per day.

We conclude that searches for the elastic scattering of LSPs have a fighting chance of success, but that this cannot be guaranteed.

5. Unstable and Metastable Relics

A number of theories predict the existence of unstable or metastable relics from the Big Bang. One example is a supersymmetric theory in which R -parity is not exact, in which case the LSP is unstable [58]. Another example is shadow matter [59], which appears in many string theories, including the original Calabi-Yau compactifications of the ten-dimensional heterotic string and [60] our own pet four-dimensional heterotic string model [61]. Shadow particles have strong non-Abelian gauge interactions among themselves, but do not join in the Standard Model gauge interactions, and so have only very weak (gravitational strength?) interactions with ordinary matter. All shadow matter bound-state particles would be very massive and unstable if the shadow sector only contained gauge supermultiplets. However, generic string models also provide matter supermultiplets in the shadow sector, which bind to form longer-lived shadow mesons and analogues of baryons [60]. Just like ordinary baryons in GUTs, these are expected to decay very slowly via higher-dimensional interactions, and so are candidates to be metastable relics. Technibaryons in some technicolour theories could behave similarly [62].

There are many astrophysical and cosmological constraints [63] on the abundance and lifetime of such an unstable or metastable relic particle, as seen in Fig. 13. One set comes from the successful confrontation of Big Bang nucleosynthesis calculations with the observed abundances of light elements. Another constraint comes from the black-body nature of the microwave background radiation, as observed by COBE in particular, which restricts relic decays that produce photons either directly or indirectly. There are also important constraints from the higher-energy cosmic γ -ray background. Finally, there are constraints from the non-observation in underground experiments of high-energy cosmic ray neutrinos from the decays of heavy relics. Figure 13 shows a compilation of all these constraints, on the basis of which we conclude that any metastable dark matter particle would need a lifetime longer than about 10^{16} y, depending on its mass.

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Fig. 11 - Regions of the $(\mu, m_{1/2})$ plane accessible to experiments detecting elastic scattering of relic neutralinos at a sensitivity of 1, 0.1, 0.01 events/kg.d on nuclei of (a) ^{73}Ge and ^{76}Ge , (b) NaI , (c) F and (d) TI [58],[59].

Fig. 12 - Plot of the fractions (50%, 25%, 12.5%) of the $(\mu, m_{1/2})$ parameter space where the relic neutralino could have the critical density in which elastic scattering off a NaI target has a rate above 0.01 events/kg.d [59].

Fig. 13 - Compilation of upper limits on the possible cosmological density of unstable or metastable relic particles [63].

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