Phenomenological Introduction to Direct Dark Matter Detection

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Abstract
The dark matter of our galactic halo may be constituted by elementary particles that interact weakly with ordinary matter (WIMPs). In spite of the very low counting rates expected for these dark matter particles to scatter off nuclei in a laboratory detector, such direct WIMP searches are possible and are experimentally carried out at present. An introduction to the theoretical ingredients entering the counting rates predictions, together with a short discussion of the major theoretical uncertainties, is here presented.

invited talk at the XXXI Rencontres de Moriond, “Dark Matter in Cosmology, Quantum Measurements, Experimental Gravitation,” Les Arcs, France, January 1996
This is a phenomenological introduction to the detection of dark matter through its scattering in a laboratory detector. For dark matter in the form of massive quasi-stellar objects, like brown dwarfs, which are much bigger and much heavier than the Earth, this type of detection is quite impracticable if not undesirable. I therefore consider dark matter in the form of elementary particles.

Many particles, most of which hypothetical, are at present candidates for dark matter: neutrinos, neutralinos, axions, etc. The methods employed in hunting for these particles are very different. In this short note I focus on this meeting’s category of particle dark matter, \textit{viz.} weakly interacting massive particles or WIMPs.

WIMPs, in a broad sense, are particles with masses of the order of atomic masses or higher \((m \approx 10\text{GeV}/c^2)\) that interact with ordinary matter with cross sections typical of the weak interaction or smaller \((\sigma \approx 10^{-38}\text{cm}^2 \text{off a proton})\). The presently most popular WIMP is the yet-undetected neutralino, the lightest supersymmetric particle in supersymmetric models. Other famous WIMPs are Dirac and Majorana neutrinos, which however, thanks to the on-going dark matter searches complemented by accelerator results, we know not to be the dominant component of our galactic halo.

A general introduction to dark matter has been given by Olive at this meeting. Direct detection of WIMPs was first explored by Goodman and Witten\(^\text{[1]}\) General reviews are Primack et al.\(^\text{[2]}\) and Smith and Lewin\(^\text{[3]}\) Engel \textit{et al.}\(^\text{[4]}\) present the nuclear physics involved. At this meeting Cabrera discusses experimental aspects of direct dark matter detection, while I focus on the theoretical aspects.

It is worth recalling some properties of the dark halo of our galaxy. Even if recent observations might change the details of our picture, the 1981 model by Caldwell and Ostriker\(^\text{[5]}\) is good for my purposes. The Sun lies at a distance of \(\approx 8.5\ \text{kpc}\) on the disk of our spiral galaxy, and moves around the center at a speed of \(\approx 220\ \text{km}/\text{s}\). The luminous disk extends to \(\approx 12\ \text{kpc}\), and is surrounded by a halo of \(\approx 100\ \text{kpc}\) where globular star clusters and rare subdwarf stars are found. Dynamical arguments suggest that the halo is filled with dark matter, whose local density in the vicinity of the Sun is estimated to be \(\rho_{\text{DM}} = 0.2–0.4\text{GeV}/c^2/\text{cm}^3 = 0.7–1.4 \times 10^{-24}\text{g/cm}^3\). Equilibrium considerations also give the root mean square velocity of halo constituents to be \(200–400\ \text{km}/\text{s}\), not much different from the escape speed from the galaxy \((500–700\ \text{km}/\text{s})\). Very little is known on the mean rotation speed of the halo, and we will assume it does not rotate. All in all, there is an optimistic factor of 2 uncertainty on the density and velocity of halo dark matter.

Are there WIMPs in our galactic halo? The scientific way to answer this question is to detect them. Signals could come from WIMP annihilation (indirect detection) or WIMP scattering (direct detection). In the former we search for rare annihilation products like neutrinos, antimatter or gamma-ray lines. This is reviewed by Bergström at this meeting. In the latter, the basic philosophy is to build a target, wait and count.

The WIMP scattering rate per target nucleus is the product of the WIMP flux \(\phi_\chi\) and of the WIMP-nucleon cross section \(\sigma_{\chi i}\). For an order of magnitude estimate we take the effective WIMP-nucleon coupling constant to be Fermi’s constant \(G_F = 2.3016 \times 10^{-19}\text{h}^2\text{c}^2/\text{cm}/\text{GeV}\), which sets the scale of weak interactions. We distinguish two cases: (i) the WIMP couples to nucleon spin, \(\sigma_{\chi i} \approx G_f^2\mu_i^2/h^4 \approx 10^{-34}\text{cm}^2\); and (ii) the WIMP couples to nucleon number, \(\sigma_{\chi i} \approx G_f^2\mu_i^2A_i/h^4 \approx 10^{-30}\text{cm}^2\). Here \(\mu_i = m_\chi m_i/(m_\chi + m_i)\) is the reduced WIMP-nucleus mass, and \(A_i\) is the atomic number of the target \((\approx 80\ \text{in the numerical examples})\). The WIMP flux is \(\phi_\chi = v\rho_\chi/m_\chi \approx 10^{-7}\text{cm}^{-2}\text{s}^{-1}/(m_\chi\text{c}^2/\text{GeV})\), for a WIMP density \(\rho_\chi \approx 10^{-24}\ \text{g/cm}^3\) and a typical WIMP velocity \(v \approx 300\ \text{km}/\text{s}\). The resulting scattering rates, taking \(m_\chi \approx 100\ \text{GeV}/c^2\), are of the order of \(\approx 1/\text{kg-day}\) for spin-coupled WIMPs and of \(\approx 10^4/\text{kg-day}\) for WIMPs coupled to nucleon number. These rates are quite small compared with normal radioactivity background.
Therefore the common denominator of direct experimental searches of WIMPs is a fight against background.

For this we get help from characteristic signatures that we do not expect for the background. For example, while the Earth revolves around the Sun, the mean speed of the WIMP “wind” varies periodically with an amplitude of 60 km/s. This leads to a ≈ 10% seasonal modulation in the detection rate, with a maximum in June and a minimum in December. As another example, the direction of the WIMP “wind” does not coincide with the Earth rotation axis, so the detection rate might present a diurnal modulation due to the diffusion of WIMPs while they cross the Earth (this however occurs for quite high cross sections). A final example of background discrimination is that the WIMP signal is directional, simply because most WIMPs come from the direction of the solar motion.

WIMP-nucleus scattering

Since the relative speed \( v \approx 300\text{km/s} \approx 10^{-3}c \), the process can be treated non-relativistically. The center of mass momentum is given in terms of the reduced WIMP-nucleus mass as \( k = \mu_i v \) and is \( \lesssim A_i \text{MeV}/c \) since \( \mu_i \lesssim m_i \). The corresponding de Broglie wavelength is \( \approx 200\text{fm}/A_i \), and can be smaller than the size of heavy target nuclei, in which case nuclear form factors are important. In the laboratory frame, the nucleus recoils with momentum \( q = 2k \sin(\theta_{\text{cm}}/2) \) and energy \( \nu = q^2/2m_i \). Here \( \theta_{\text{cm}} \) is the center-of-mass scattering angle. The 4-momentum transfer is very small, \( Q^2 < A_i^2 10^{-6}\text{GeV}^2/c^2 \) (compare with a typical deep inelastic \( Q^2 > 1\text{GeV}^2/c^2 \)).

The differential scattering rate per unit recoil energy and unit target mass is formally

\[
\frac{dR}{d\nu} = \frac{\rho_\chi}{m_\chi} \sum_i f_i \eta_i(q) \sqrt{\frac{|T_i(q^2)|^2}{2\pi\hbar^4}}.
\]

The sum is over the nuclear isotopes in the target, \( T_i(q^2) \) is the scattering matrix element at momentum transfer squared \( q^2 = 2m_i\nu \), and \( f_i \) is the mass fraction of isotope \( i \). A sum over final and average over initial polarizations is understood in \( |T_i(q^2)|^2 \). The factor

\[
\eta_i(q) = \int^{\infty}_{q/2\mu_i} \frac{f_\chi(v)}{v} dv
\]

with units of inverse velocity, incorporates the \( \chi \) velocity distribution \( f_\chi(v) \). For a Maxwellian distribution with velocity dispersion \( v_{\text{rms}} \), seen by an observer moving at speed \( v_O \),

\[
\eta_i(q) = \frac{1}{2v_O} \left[ \text{erf} \left( \frac{v_q + v_O}{\sqrt{2}v_{\text{rms}}} \right) - \text{erf} \left( \frac{v_q - v_O}{\sqrt{2}v_{\text{rms}}} \right) \right],
\]

with \( v_q = q/2\mu_i \). For standard halo parameters, \( \eta_i(q) \) is approximately exponential in the deposited energy \( \nu \). The previously-mentioned modulations enter the rate through \( \eta_i(q) \).

The scattering matrix element \( T(q^2) \) can be written as the Fourier transform

\[
T(q^2) = \int \langle f|V(\vec{r})|i \rangle e^{i\vec{q}\cdot\vec{r}}/\hbar d\vec{r}
\]

of a non-relativistic WIMP-nucleus potential

\[
V(\vec{r}) = \sum_{n=p,n} \left( G_n^p + G_n^\sigma \vec{\sigma} \vec{\sigma}_n \right) \delta(\vec{r} - \vec{r}_n).
\]
The constants $G^n_s$ and $G^n_a$ are effective four-fermion coupling constants for nucleon-WIMP interactions, and are analogous to Fermi’s constant $G_F$. $G^n_s$ represents scalar[1] or spin-independent interactions, $G^n_a$ axial[1] or spin-dependent interactions. Both terms are coherent in the quantum-mechanical sense when $qR_{\text{nucleus}} \ll \hbar$, i.e. when the nucleus can be treated as pointlike and $T(q^2)$ can be taken as $T(0)$. At larger $q$, which can occur with heavy target nuclei, both terms are incoherent. Nuclear form factors $F(q^2)$, conventionally defined by $T(q^2) = T(0)F(q^2)$, should then be introduced. The scalar and spin form factors are in general different, reflecting the difference in the mass and spin distributions inside the nucleus.

The task of a theoretician is to provide a theoretical estimate of $T(q^2)$ starting from a particle-physics model. We accomplish this by stages, successively finding the WIMP-quark, the WIMP-nucleon and the WIMP-nucleus effective lagrangians. Step 1, finding the effective WIMP-quark lagrangian at small $q^2$, is analogous to going from the Standard Model to four-fermion interactions. Step 2 requires knowledge of the quark content of the nucleon, i.e. the contributions of different quarks to the nucleon mass and spin. Step 3 needs a nuclear model to describe how protons and neutrons are distributed in a nucleus.

This procedure is now illustrated for a Dirac neutrino and for a Majorana particle, an example of which is the neutralino.

**Dirac neutrino**

Step 1: a Dirac neutrino $\nu$ interacts with a quark $q$ through the diagram in Fig. 1a. At $q^2 \ll m_Z^2$, the $Z$ propagator reduces to $ig \gamma^\mu/m_Z^2$, and the four-fermion amplitude reads

$$\sqrt{2}G_F \bar{\nu}(v_\nu - a_\nu \gamma_5)\gamma^\mu \nu (v_q - a_q \gamma_5)\gamma^\mu q,$$

with $v_\nu = a_\nu = \frac{1}{2}$, $a_q = T_{3q}$ and $v_q = T_{3q} - 2e_q \sin^2 \theta_W$. Here $\sin^2 \theta_W \simeq 0.23$ and $e_q$ and $T_{3q}$ are the electric charge and the third component of the weak isospin of quark $q$. For a non-relativistic neutrino, only the time component of the vector current and the space components of the axial current survive. The first is spin-independent ($\bar{\nu}\gamma^0\nu \propto \bar{\nu}v\nu$) and the second spin-dependent ($\bar{\nu}\gamma^5\nu \propto \bar{\nu}\sigma\nu$).

Step 2 for the vector part

$$\sqrt{2}G_F v_\nu v_q \bar{\nu}\gamma^\mu \nu \bar{q}\gamma^\mu q,$$

leads to the four-fermion coupling constants

$$G^p_s = \frac{G_F}{\sqrt{2}}(1 - 4 \sin^2 \theta_W)v_\nu$$

$$G^n_s = -\frac{G_F}{\sqrt{2}}v_\nu.$$  

The interaction is mainly with the neutrons since $1 - 4 \sin^2 \theta_W \approx 0$.

Step 2 for the axial part

$$\sqrt{2}G_F a_\nu a_q \bar{\nu}\gamma^\mu \gamma_5 \nu \bar{q}\gamma^\mu q,$$

leads to the four-fermion coupling constants

$$G^p_a = \sqrt{2}G_F a_\nu (a_u \Delta u + a_d \Delta d + a_s \Delta s),$$

$$G^n_a = \sqrt{2}G_F a_\nu (a_u \Delta d + a_d \Delta u + a_s \Delta s).$$

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1 Associated to scalar and axial vectors under 3d rotations.
Figure 1: Examples of WIMP–quark scattering.

Here $\Delta q$ is the fraction of the proton spin carried by quark $q$. $\frac{1}{2}\langle p|\bar{q}\gamma^5\gamma_5|p\rangle = \Delta q s_\mu$. It can be obtained from data on neutron and hyperon $\beta$-decay, which give $\Delta u - \Delta d = 1.2573 \pm 0.0028$ and $\Delta u + \Delta d - 2\Delta s = 0.59 \pm 0.03$, respectively. The contribution of the strange quark is $\Delta s = 0$ in the naive quark model, $\Delta s = -0.11 \pm 0.03 \pm \cdots$ from deep inelastic data, and $\Delta s = -0.15 \pm 0.09$ from elastic $\nu p \rightarrow \nu p$ data.

Step 3 for the spin-independent part introduces the nuclear mass form factor $F_{\text{mass}}(q^2)$, and results in

$$|T_a(q^2)|^2 = |ZG_p^a + NG_n^a|^2 |F_{\text{mass}}(q^2)|^2,$$

where $N$ ($Z$) is the number of neutrons (protons) in the nucleus. Neutron scattering off nuclei suggests that $F_{\text{mass}}(q^2) \approx F_{\text{e.m.}}(q^2)$, the electromagnetic form factor. The electric charge distribution is well-described by a Fermi or Woods-Saxon form, whose Fourier transform is indistinguishable from the convenient analytic expression

$$F_{\text{e.m.}}(q^2) \approx \frac{3j_1(qR)}{qR} e^{-\frac{1}{2}(qs)^2}.$$

The electromagnetic radius $R$ and the surface thickness $s$ can be obtained by fitting electron scattering data or can be roughly approximated by $R \approx A^{1/3}$ fm and $s \approx 1$ fm. $F_{\text{e.m.}}(q^2)$ presents diffraction zeros when the modified Bessel function $j_1(qR) = 0$, the first of which occurs at $qR \approx 4.2$. In electron scattering, these diffraction zeros are filled in, because due to the long-range Coulomb attraction the electron wave function is distorted from a simple plane wave and the form factor is not simply the Fourier transform of the charge density. The short-range nature of WIMP-nucleus interactions make us expect no wave function distortion, and diffraction zeros remain. The first diffraction zero is important in assessing bounds from some present-day detectors.

Step 3 for the spin-dependent part requires the expectation values of the total spin of protons $\langle S_p \rangle$ and neutrons $\langle S_n \rangle$ separately. At $q = 0$,

$$|T_a(0)|^2 = \frac{4(J + 1)}{J} |G_p^a \langle S_p \rangle + G_n^a \langle S_n \rangle|^2,$$

where $J$ is the nuclear spin. Even-even nuclei, with even numbers of protons and of neutrons, do not have spin, and for them $T_a(0) = 0$. For even-odd nuclei with $J \neq 0$, a nuclear model is

$^2$The first corrections are at a level of $10^{-6}$ and come from neglected higher powers of the incoming WIMP velocity.
needed to estimate $\langle S_p \rangle$ and $\langle S_n \rangle$. For instance, $^{73}\text{Ge}$ is an odd-neutron nucleus with $J = \frac{9}{2}$. The single-particle shell model\cite{1,13} gives

$$\langle S_n \rangle = \frac{1}{2} \left[ 1 + \frac{3}{4} - \frac{l(l+1)}{j(j+1)} \right] = 0.50, \quad \langle S_p \rangle = 0;$$

(16)

the odd-group model\cite{14} in which the odd-nucleon spin is related to the nuclear magnetic moment $\mu$ and gyromagnetic factors $g_{L,S}^n$, gives

$$\langle S_n \rangle = \frac{\mu - g_{S}^L J}{g_{S}^n - g_{L}^n} = 0.23, \quad \langle S_p \rangle = 0;$$

(17)

a more sophisticated interacting shell model\cite{15,16} gives

$$\langle S_n \rangle = 0.468, \quad \langle S_p \rangle = 0.011.$$  

(18)

The proton might have a small but non-zero contribution to the cross section, which might change the relative merits of different nuclei for dark matter searches.

At $q \not= 0$, nuclear spin form factors are needed. The neutron and proton contributions differ, and at present only complex calculations\cite{15,16} for specific nuclei provide an estimate of the isoscalar and isovector spin form factors $F_{\text{spin}}^0(q^2)$ and $F_{\text{spin}}^1(q^2)$, in terms of which

$$|T_a(q^2)|^2 = \frac{J+1}{J} \left| (G_p^p + G_n^n) \langle S_p + S_n \rangle F_{\text{spin}}^0(q^2) + (G_p^p - G_n^n) \langle S_p - S_n \rangle F_{\text{spin}}^1(q^2) \right|^2,$$

(19)

The results of these calculations can be conveniently resumed by the approximate expressions

$$F_{\text{spin}}^0(q^2) \simeq \exp \left( -\frac{r_0^2 q^2}{\hbar^2} \right), \quad F_{\text{spin}}^1(q^2) \simeq \exp \left( -\frac{r_1^2 q^2}{\hbar^2} + i \frac{c q}{\hbar} \right),$$

(20)

with parameters given in the following table for selected nuclei:

| $J$  | $\nu_{\text{max}}$/keV | $\langle S_p \rangle$ | $\langle S_n \rangle$ | $r_0$/fm | $r_1$/fm | $c$/fm | $\text{valid for } \nu/\text{keV} <$ |
|------|-----------------|-----------------|-----------------|---------|---------|--------|------------------|
| $^{73}\text{Ge}$ | $\frac{9}{2}$ | 540 | 0.011 | 0.468 | 1.971 | 2.146 | -0.246 | 55 |
| $^{28}\text{Si}$ | $\frac{1}{2}$ | 216 | -0.0019 | 0.133 | 1.302 | 1.548 | -0.320 | 145 |
| $^{27}\text{A}$ | $\frac{1}{2}$ | 100 | 0.3430 | 0.269 | 1.378 | 1.600 | 0.196 | $\nu_{\text{max}}$ |
| $^{39}\text{K}$ | $\frac{1}{2}$ | 145 | -0.184 | 0.054 | 1.746 | 1.847 | 0.371 | $\nu_{\text{max}}$ |

**Majorana fermion**

A Majorana fermion is a spin-$\frac{1}{2}$ particle that coincides with its antiparticle. It carries no conserved quantum number. It has neither vector nor tensor currents. Of the remaining pseudoscalar, scalar and axial currents, only the last two have a non-vanishing non-relativistic limit, spin-independent the first ($\bar{\chi}\chi \propto \chi^\dagger \chi$) and spin-dependent the second ($\bar{\chi}\gamma_5 \chi \propto \chi^\dagger \vec{\sigma} \chi$).

Axial currents may arise from exchange of a Z boson as in fig. 1b, and the analysis is then analogous to that in the previous section, with the obvious replacement of $a_\nu$ with $a_\chi$.

Scalar currents originate from exchange of a scalar particle $\varphi$, e.g. as in Fig. 1c. At small $q^2$, the $\varphi$ propagator reduces to $-i/m_\varphi^2$ and the four-fermion amplitude reads

$$- \frac{g_{\varphi\chi\chi} g_{\varphi\varphi}}{m_\varphi^2} \bar{\chi}\chi \bar{\varphi}\varphi.$$  

(21)
Figure 2: Scattering rate versus mass for neutralinos: (a) phenomenological approach, (b) grand-unified approach.

For a nucleon \( n = p, n \) one then obtains
\[
G^n_s = -\frac{g_\phi \chi \chi}{m^2_\phi} \sum_q g_{\phi q q} \langle n|\bar{q} q|n\rangle.
\]
(22)

For example, in the case of the neutralino with exchange of the lightest supersymmetric Higgs boson, the sum over quarks is explicitly
\[
\frac{g}{2m_W} \left[ \frac{\cos \alpha}{\sin \beta} \langle m_u \bar{u} u + m_c \bar{c} c + m_t \bar{t} t \rangle - \frac{\sin \alpha}{\cos \beta} \langle m_d \bar{d} d + m_s \bar{s} s + m_b \bar{b} b \rangle \right].
\]
(23)

The scalar quark content of the nucleon \( \langle n|\bar{q} q|n\rangle \) can be extracted from data with the help of chiral perturbation theory, \( \pi \)-nucleon scattering and heavy quark expansion. The result is
\[
\langle m_u \bar{u} u \rangle \simeq \langle m_d \bar{d} d \rangle \simeq 30 \text{ MeV}/c^2, \quad \langle m_s \bar{s} s \rangle \simeq 60–120 \text{ MeV}/c^2,
\]
(24)
\[
\langle m_c \bar{c} c \rangle = \langle m_b \bar{b} b \rangle = \langle m_t \bar{t} t \rangle = \frac{2}{27} \left( m_p - \sum_{q=u,d,s} \langle m_q \bar{q} q \rangle \right) \simeq 60 \text{ MeV}/c^2.
\]
(25)

The strange quark contribution is uncertain by a factor of 2. Step 3 is analogous to the Dirac neutrino case, and leads to eq. (13) with four-fermion couplings given by (22).

**Neutralino**

Supersymmetry and the neutralino have been presented by Jungman at this conference. The neutralino has both spin-dependent and spin-independent interactions with nuclei, the former mediated by Z boson and squarks, the latter by Higgs bosons and squarks. The general formalism of the preceding sections can be used. In the limit of heavy squarks \( \tilde{q}_k \), the effective four-fermion constants are given by
\[
G^p_s \simeq G^n_s = \sum_{q=u,d,s,c,b,t} \langle \bar{q} q \rangle \left( - \sum_{h=H_1,H_2} \frac{g_{h \chi \chi} g_{h q q}}{m^2_h} + \frac{1}{2} \sum_{k=1}^{6} \frac{g_{L \tilde{q}_k \chi q} g_{R \tilde{q}_k \chi q}}{m^2_{\tilde{q}_k}} \right),
\]
(26)
Expressions for the elementary vertices \( g_{ijk} \) can be found in ref. 18.

Predictions in supersymmetric models suffer from the presence of many unknown parameters. Two extreme attitudes are a phenomenological approach in which what is not excluded is allowed, and a grand-unified approach in which coupling constants and masses are unified at some high energy scale. Fig. 2 shows examples of calculated event rates in \(^{76}\text{Ge}\), each point representing a choice of model parameters: “predictions” may well span 10 orders of magnitude in a phenomenological approach\(^{18}\) and 2 orders of magnitude in a more restricted scenario\(^{19}\).

**Underabundant dark matter relics**

Given a particle-physics model, the relic density of a species, a WIMP \( \chi \) in particular, is a calculable and definite quantity. Often it happens that the computed relic density \( \Omega_\chi \) is (much) smaller than the dark matter density in the Universe. For this reason, some authors simply neglect this case. But even if these WIMPs constitute only a fraction of the dark matter, they generally have quite high scattering cross sections off nuclei, because of an approximate inverse proportionality of the \( \chi \) relic density and the \( \chi \)–nucleus cross section. However, the scattering rate also includes the \( \chi \) halo density \( \rho_\chi \). It is reasonable that \( \rho_\chi \) is only a fraction of the local dark matter density \( \rho_{\text{DM}} \), but which precise fraction it is depends on the model for galaxy formation. If both the main and the \( \chi \) components of dark matter are cold, we expect them to behave similarly under gravitation, so that the halo fraction \( f_\chi \) might be equal to the universal fraction \( \Omega_\chi/\Omega_{\text{DM}} \). Unfortunately, \( \Omega_{\text{DM}} \) is poorly known: it can range from \( \approx 0.01 \) for dark matter associated with galactic halos to \( \approx 1 \) for a smooth universal component. In fig. 3, the suppression of scattering rates due to rescaling of the neutralino halo density by a universal fraction with \( \Omega_{\text{DM}} h^2 = 0.025 \) is apparent to the left of the dashed line. This suppression must be included for consistency when setting bounds on particle-physics models.

![Figure 3: Scattering rate versus relic density for neutralinos (from ref. 18).](image)

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