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INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

CHIRALITY, PHASE TRANSITIONS AND THEIR INDUCTION IN AMINO ACIDS

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INTERNATIONAL ATOMIC ENERGY AGENCY

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ABSTRACT

"Atoms such as carbon, oxygen, nitrogen and hydrogen, the major constituents of biological molecules, are less than 0.4 nm in diameter.... The behaviour of small molecules is a reflection of the intrinsic properties of the constituent atoms. Hence it might be expected that the behaviour of large macromolecules can be explained by a knowledge of atomic properties. Since organelles, whole cells and organisms are essentially macromolecular assemblies, it may be possible in time to derive an atomic theory of life'' (Rees and Sternberg) [1].

It has been suggested that chirality among the twenty amino acids which make up the proteins may be a consequence of a phase transition which is analogous to that due to BCS superconductivity [2]. We explore these ideas in this paper and show, following Lee and Drell [3], that a crucial form for the transition temperature T_c involves dynamical symmetry breaking. The t-quarks or supersymmetry (or something similar which ensures a heavy mass) appear to be essential if such mechanisms are to hold.

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1. The T_c for BCS superconductivity for metals is of the form $\omega \exp(-\frac{2}{g_{eff}\sigma(0)})$. We conjectured that a similar formula may hold for the case of amino acids chains. In the present paper we shall explore this further.

First, let us review the subject of BCS superconductivity. The best treatment, which I know of has been given by Sakita [4], following on the work of Ginsberg and Landau [5] and of Gorkov [6].

The idea is to start from the Feynman Lagrangian methodology of writing down the BCS theory for the superconducting electronic system. One tries to write down the equivalent Ginzburg-Landau equation. From this equation is deduced the value of temperature T_c.

2. The BCS *Hamiltonian* is given by

$$\begin{aligned} H &= H_0 + H_1 \\ H_0 &= \sum_{s=\uparrow\downarrow} \int d\vec{x}\psi_s^\dagger(\vec{x})(-\nabla\frac{2}{2m} - \mu)\psi_s(\vec{x}) \\ H_1 &= -g_{eff}\int d\vec{x}\psi_\uparrow^\dagger(\vec{x})\psi_\downarrow^\dagger(\vec{x})\psi_\downarrow(\vec{x})\psi_\uparrow(\vec{x}) \end{aligned} \tag{2.1}$$

where g_{eff} is an attractive coupling constant between spin up ($\uparrow$) and spin down ($\downarrow$) electrons and antielectrons is the chemical potential. *The sign of g_{eff} is part of the assumption* of the Hamiltonian which signifies an attractive force between Cooper-paired systems of electrons consist of one of the particles being replaced by its antiparticle with a factor of two which appears in the mass term.

We can introduce electro-magnetic interaction in a gauge-invariant way by the minimal substitution

$$\vec{\nabla} \to \vec{\nabla} - ie\vec{A}$$

and treat the vector potential $\vec{A}(\vec{x})$ as an external source. Then, in general, the Hamiltonian equals

$$\begin{aligned} H = \sum_{s=\uparrow\downarrow}\int d\vec{x}\psi_s^\dagger(\vec{x})[-\frac{1}{2m}(\vec{\nabla} - ie\vec{A})^2 - \mu]\psi_s(\vec{x}) \\ -g_{eff}\int d\vec{x}\psi_\uparrow^\dagger(\vec{x})\psi_\downarrow^\dagger(\vec{x})\psi_\downarrow(\vec{x})\psi_\uparrow(\vec{x}) \end{aligned} \tag{2.2}$$

which is invariant under gauge transformations: –

$$\psi_s(\vec{x}) \to e^{ie\Lambda(\vec{x})}\psi_s(\vec{x}), \quad \vec{A}(\vec{x}) \to \vec{A}(\vec{x}) + \vec{\nabla}\Lambda(\vec{x}) \,. \tag{2.3}$$

The partition function based on the appropriate Lagrangian corresponding to the above Hamiltonian is given by

$$Z = \int D\psi D\bar{\psi} \exp\left[-\int_0^\beta d\tau \int_V d\vec{x}\bar{\psi}_s\left(\frac{\partial}{\partial\tau} - \frac{1}{2m}(\vec{\nabla} - ie\vec{A})^2 - \mu\right)\psi_s\right]$$

where we have written a non-relativistic equation for the fermion (electron) ψ in the theory.

$$\exp\left[g_{eff}\int_0^\beta d\tau\int_V d\vec{x}\bar{\psi}_\uparrow(\vec{x})\bar{\psi}_\downarrow(\vec{x})\psi_\downarrow(\vec{x})\psi_\uparrow(\vec{x})\right] \tag{2.4}$$

Note that the sign before g_{eff} has changed to plus g_{eff} instead of minus g_{eff} in the Hamiltonian formulation. This is because $\Delta H = -\Delta L$ for the case of potentials which do not contain time derivatives of the fields themselves. The 4-fermion interaction can be expressed in terms of a complex auxiliary Higgs scalar field ϕ:

$$\frac{1}{C}\int D\phi D\phi^* \exp[+\kappa^2\int d^4x\phi^*\phi + g_{eff}^{1/2}\kappa\int d^4x(\bar{\psi}_\uparrow\bar{\psi}_\downarrow\phi + \psi_\downarrow\psi_\uparrow\phi^*)] \tag{2.5}$$

where we have used the notation

$$\int d^4x \equiv \int_0^\beta d\tau\int_V d^3\vec{x} \tag{2.6}$$

The constant C is given by

$$C = \int D\phi D\phi^* \exp[+\kappa^2\int d^4x\phi^*\phi] \tag{2.7}$$

where κ is a constant with dimension of mass. We shall introduce a source for the ϕ field in the form

$$Z[j,j^*] = \frac{1}{C}\int D\psi D\bar{\psi}D\phi D\phi^* \exp\left[-\int d^4x\bar{\psi}_s\left(\frac{\partial}{\partial\tau} - \frac{(\vec{\nabla}-ie\vec{A})^2}{2m} - \mu\right)\psi_s\right]$$

$$\times \exp[+\kappa^2\int\phi^*\phi d^4x + g_{eff}^{1/2}\kappa\int d^4x(\bar{\psi}_\uparrow\bar{\psi}_\downarrow\phi + \psi_\downarrow\psi_\uparrow\phi^* + j^*\phi + \phi^*j)] \tag{2.8}$$

The partition function is given by

$$Z = Z[j,j^*]|_{j=j^*=0} \tag{2.9}$$

The action which has appeared in (2.8) is $Ld\tau$ where

$$L = L_0 + L_1 \tag{2.10}$$

$$L_0 = \int d\vec{x}\left[\bar{\psi}_s\left(\frac{\partial}{\partial\tau} - \frac{\vec{\nabla}^2}{2m} - \mu\right)\psi_s + \kappa^2\phi^*\phi\right] \tag{2.11}$$

$$L_1 = \int d\vec{x}\Big[-\frac{ie}{2m}(\bar{\psi}_s\vec{\nabla}\psi_s - (\vec{\nabla}\bar{\psi}_s)\psi_s)\vec{A} + \frac{e^2}{2m}\bar{\psi}_s\psi_s\vec{A}^2 +$$

$$+ g_{eff}^{1/2}\kappa(\bar{\psi}_\uparrow\bar{\psi}_\downarrow\phi + \psi_\downarrow\psi^\uparrow\phi^*)\Big] \tag{2.12}$$

Using the temperature dependent Fourier decomposition

$$\psi_s(\vec{x},\tau) = (\beta V)^{-1/2}\sum_k \psi_s(\vec{k},n)e^{i(\vec{k}\cdot\vec{x}-\xi_n\tau)} \tag{2.13}$$

we obtain the free action

$$\int_0^\beta d\tau L_0 = \sum_{\vec{k},n}\left[\bar{\psi}_s(\vec{k},n)(-i\xi_n+\omega_k)\bar{\psi}_s(\vec{k},n) + \kappa^2\bar{\phi}^*(\vec{k},n)\bar{\phi}(\vec{k},n)\right] \tag{2.14}$$

Therefore, the propagator of the electron with four momentum $(\vec{k},n)$ is given by

$$\frac{1}{\omega_k - i\xi_n}, \quad \text{where} \quad \omega_k = \frac{\vec{k}^2}{2m} - \mu, \quad \xi_n = \frac{\pi}{\beta}(2n+1) \tag{2.15}$$

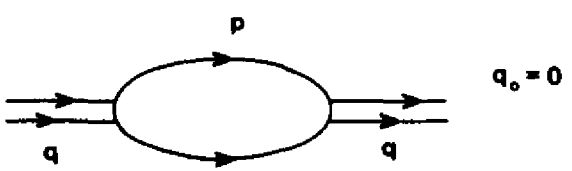
and the propagator of ϕ by

$$\frac{1}{\kappa^2} \tag{2.16}$$

Notice that the propagator of ϕ to this order does not depend on four momentum, because the Lagrangian (2.11) does not contain the kinetic energy term.

3. Calculation of the loop diagram:

We shall compute 1-loop diagram with 2 external ϕ-lines. The (temperature dependent) leading terms are



$$= \frac{g_{eff}\kappa^2}{\beta V}\sum_{n,\vec{p}}\frac{1}{(i\xi_n-\omega_p)(-i\xi_n-\omega_{p'})}$$

$$= A + B\vec{q}^2 + \ldots \tag{3.1}$$

where

$$A = \frac{g_{eff}\kappa^2}{\beta V}\sum_{n,\vec{p}}\frac{1}{\xi_n^2+\omega_p^2} = \sum_n\frac{g_{eff}\kappa^2}{\beta}\frac{1}{(2\pi)^3}\int d^3\vec{p}\frac{1}{\xi_n^2+\omega_p^2} \tag{3.2}$$

Introducing density of states $\sigma(\omega)$ from the ansatz

$$\frac{1}{(2\pi)^3}\int d^3\vec{p} \equiv \int d\omega_p\sigma(\omega_p) \tag{3.3}$$

we get

$$A = \sum_n \frac{g_{eff}\kappa^2}{\beta} \int_{-\infty}^{\infty} \frac{d\omega\sigma(\omega)}{\xi_n^2 + \omega^2} \approx \sum_n \frac{g_{eff}\kappa^2}{\beta}\sigma(0) \int_{-\infty}^{\infty} \frac{d\omega}{\xi_n^2 + \omega^2} \tag{3.4}$$

where $\omega_p = \frac{p^2}{2m} - \mu = 0$ is the Fermi surface and we have approximated by using density of states at the Fermi surface by putting a cut-off n_{max}

$$\omega_D = k_B T \pi (2n_{max} + 1) \tag{3.5}$$

Thus

$$A = g_{eff}\kappa^2 \sigma(0) \ln \frac{4\gamma\omega_D}{\pi k_B T}, \quad \text{where } \gamma = \text{Euler constant} \tag{3.6}$$

The coefficient of $|\phi_c|^2$ terms $\Gamma[\phi_c, \phi_c^*]$ is then

$$\kappa^2 \left(1 - g_{eff}\sigma(0) \ln \frac{4\gamma\omega_D}{\pi k_B T}\right) \phi_c^* \phi_c \tag{3.7}$$

The expression within the brackets gives the critical temperature as

$$T_c = \frac{4\gamma\omega_D}{\pi k_B} \exp\left(-\frac{1}{g_{eff}\sigma(0)}\right) \tag{3.8}$$

Note that the terms appearing before the logarithm in (3.7) appear in the exponent while terms which are arguments for the log appear as $A_D\omega$ in the formula for T_c.

Since the Higgs mechanism is the relativistic version of the above, one may write the entire calculation of one-loop diagrams as well as the calculation of T_c in the form of corrections to the calculation of mass of the Higgs field. We shall adopt this procedure with slight variations.

4. The point we wish to make is summarized by (3.8) which has the form $A\omega_D \exp\left(-\frac{2}{g_{eff}\sigma(0)}\right)$. In order that exp factor does not present too much of a restriction $\frac{1}{g_{eff}\sigma(0)}$ should be of the order of unity. This gives $T_c \approx \frac{\gamma\omega_D}{\pi k_B} \times \exp -2$.

Before proceeding, we shall make some remarks about the amino acids. The crystalline structure of amino acids is characterized by the graph shown in figure 1.

Apart from the N terms, here R_1 to R_4 are the residues which specify the amino acid concerned. The peptide bond which gives rise to proteins is formed, for example, where the molecules of water are consistently expelled and the lattice structure is reduced to the simple form

The hydrogens give up their loose electrons and act as metallic hydrogens. (Superconductivity has been established for a similar case by W.J. Carr [7]).

5. Helicity and chirality of the matter fields.*

Let us review the definition of chirality, which is the eigenvalue of γ_5, with $\gamma_5 = +1$ corresponding to right-handedness, and $\gamma_5 = -1$ to left-handedness:

$$\gamma_5 R = R, \quad \bar{R}\gamma_5 = -\bar{R},$$

$$\gamma_5 L = -L, \quad \bar{L}\gamma_5 = \bar{L}, \tag{5.1}$$

where R, L are Dirac spinors with only two independent components. They may be obtained from a 4-component Dirac spinor ψ by the following projections:

$$R = \frac{1}{2}(1 + \gamma_5)\psi, \quad \bar{R} = \frac{1}{2}\bar{\psi}(1 - \gamma_5),$$

$$L = \frac{1}{2}(1 - \gamma_5)\psi, \quad \bar{L} = \frac{1}{2}\bar{\psi}(1 + \gamma_5). \tag{5.2}$$

For later applications, it is important to note that

$$\bar{\psi}\psi = \bar{L}R + \bar{R}L,$$

$$\bar{\psi}\gamma^\mu\psi = \bar{L}\gamma^\mu L + \bar{R}\gamma^\mu R. \tag{5.3}$$

written as

$$(\alpha \cdot \mathbf{p} + \beta m)\psi = E\psi, \quad E = (p^2 + m^2)^{1/2}. \tag{5.4}$$

Using the identity $\alpha = \gamma_5\sigma$, and the fact that γ_5 and σ commute, we can rewrite this in the form

$$\sigma \cdot \hat{\mathbf{p}} R = \frac{E}{p}R - \frac{m}{p}\beta L, \quad p = |\mathbf{p}|,$$

$$\sigma \cdot \hat{\mathbf{p}} L = -\frac{E}{p}L - \frac{m}{p}\beta R. \tag{5.5}$$

These equations become decoupled if $m = 0$:

$$\sigma \cdot \hat{\mathbf{p}} R = R$$

$$\sigma \cdot \hat{\mathbf{p}} L = -L. \tag{5.6}$$

* This chapter is taken essentially from the book "Quarks, Leptons and Gauge Fields" Kerson Huang (World Scientific, 1982).

Therefore, for massless Dirac particles, chirality is the same as helicity, for antiparticles, chirality is the opposite of helicity. [An antiparticle has the same chirality as the particle, by definition; but it has the opposite helicity due to a change in the sign of E in (5.4)].

A conventional mechanical mass term in the Lagrangian density cannot be invariant under $SU(2)$, because it is proportional to $\bar{\psi}\psi = \bar{L}R + \bar{R}L$. Therefore, in this theory the electron mass can arise only by virtue of a spontaneous breakdown of SU(2). A convenient way to implement this is to introduce a doublet Higgs field

$$\phi = \begin{pmatrix} \phi_+ \\ \phi_0 \end{pmatrix} , \tag{5.7}$$

where the subscripts refer to the electric charges, write the mass term as

$$\mathcal{L}^{\text{mass}} = \vartheta(\bar{L}\phi R + \bar{R}\phi^\dagger L)(\frac{\rho}{\rho_0}) , \tag{5.8}$$

where $\bar{L}\phi$ is an $SU(2)$ singlet, and a Dirac spinor. If ϕ has non-zero vecuum value, then for low excitations and $\rho = \rho_0$, (5.8) is indistinguishable from a conventional mass term.

The Weinberg-Salam model is obtained by gauging $SU(2) \times U(1)$, generated by weak hypercharge.

With these, the covariant derivative can be written in the form

$$D^\mu = \partial^\mu + ig(W_1^\mu t_1 + W_2^\mu t_2) + ieQA^\mu + ieQ'Z_\mu , \tag{5.9}$$

Electric charge $= eQ$, $\quad Q = \tau_3 + \tau_0$ where the neutral charge matrix Q' is defined by

$$eQ' = \tau_3 \cot\theta_w - \tau_0 \tan\theta_w . \tag{5.10}$$

Here τ_0 is the singlet matrix among $\tau_0, \tau_1, \tau_2, \tau_3$, while $Q = \tau_3 + \tau_0$

To study the masses of the gauge fields, it is convenient to go to unitary gauge, in which

$$\phi = \begin{pmatrix} 0 \\ \frac{\rho}{\rho_0} \end{pmatrix} . \tag{5.11}$$

where $\frac{\rho}{\rho_0}$ is a real field.

We should in fact have used a triplet of fields $\begin{pmatrix} W^+ \\ Z^0 \\ W^- \end{pmatrix}$ $W^+ = W_1 - \frac{iW_2}{\sqrt{2}}$ likewise $W^- = W_1 + \frac{iW_2}{\sqrt{2}}$ belonging to a vector representation of the internal group SU(2), while the Higgs particles are a complex $\frac{1}{2}$ representation of the same group SU(2). We have neglected W^+ and W^- throughout this paper because their contribution is small at low momenta.

In terms of fields in the unitary gauge, the Lagrangian density is

$$\mathcal{L} = -\frac{1}{4}(\mathbf{G}\cdot\mathbf{G} + \mathbf{H}\cdot\mathbf{H}) + \frac{1}{4}g^2\rho^2\left(W_1^2 + W_2^2 + \frac{Z^2}{\cos^2\theta_w}\right) + \partial\rho\cdot\partial\rho - \lambda(\rho^2 - \rho_0^2)^2 , \tag{5.12}$$

Rewrite in the unitary gauge, where η is the real Higgs field in unitary gauge, in which $\phi(x)$ has the form (5.12) and

$$\rho(x) = \rho_0 + \eta(x) . \tag{5.13}$$

6. To compute any further, we must guarantee the g_{eff} in the effective Lagrangian is a positive number corresponding to attraction. Further, the proton left behind migrates to the nitrogen so that the Z^0 meson before being annihilated gives rise to an interaction between the quarks contained inside the proton and inside the nitrogen. (That it has to be nitrogen rather than the carbon (next door) is empirically guaranteed by the "right" configuration which is imparted to the corresponding molecule which contains sugars * and no nitrogens). **

To consider Z^0 containing part of the Lagrangian, write it as

$$L_{int}(Z^0) = \frac{eZ^0}{\sin\theta\cos\theta}[(T_{3L} - \sin^2\theta J_{em\ L+R})] . \tag{6.1}$$

Here J_{em} is the electromagnetic current, T_{3L} is the left-handed third component of the weak isospin consisting of the (anomaly-free) combination of the proton and neutron (p, n), (or the quarks inside

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* We may consider 2,8,20,28, 50,82 and 126 as "magic numbers" for nuclei, provided spin-orbit coupling is taken into account. (See paper by M. Goeppert Mayer)[8]). Nuclei containing 2,8,20,28,50,82 or 126 neutrons or protons are particularly stable. The detailed evidence supporting this point of view is discussed in Ref. 9 with the fact of 20 coming out naturally. M. Goeppert Mayer goes on to consider in detail the protons and neutrons and the spin-orbit couplings in terms of a potential energy which has a shape somewhat between that of a square well and a three-dimensional isotropic oscillator. (See Table 1).

** The discussion from now on is not exclusively relevant to the rest of this paper. The problem for the usual quark model is to see if the Pauli principle holds for quark assignments. This means that in the final analysis the quark model which substitutes quarks for nucleons will have quark spheres of influence reduced depending on the number of colours. This implies that certain discrepancies in the nucleon model are removed when quarks are taken into account. These discrepancies, for example, concern the spin of ${}_{11}Na^{12}$. The magnetic moment of this nucleus would indicate $p_{3/2}$ rather than $d_{3/2}$ orbit. That these considerations have something new to tell us is an important point in its favour. (We would urge very strongly that parity violation among the proteins used by this mechanism should be further examined in order that the Z^0 and its decay are properly considered).

the proton or the neutron). Z^0 here requires a mass because of the sponeaneous symmetry breaking implied by the last equation.

$$L^{\text{mass}} = g_{eff} \times (T_L^3 - (4 \sin^2 \theta J_{em\ L+R})) \times L_{int} \times \frac{\rho}{\rho_0} \qquad (6.2)$$

Here g_{eff} is proportional to $\frac{(e^2)}{16 \sin^2 \theta \cos^2 \theta} \times \frac{\eta_{\mu\nu}}{k^2 - m_Z^2}$ where $\eta_{\mu\nu} = 1, -1, -1, -1$. This is positive provided $m_Z^2 \gg k^2$ and we consider the 3 space-rotation invariant part of the expression above.

We shall take $(1 - 4 \sin^2 \theta) \approx \frac{1}{13}$ with the present *empirical value* of the parameter $\sin^2 \theta \approx .231$. We shall not take this quantity to equal zero as has been done by authors of Refs. 9 and 10 in the hope that the renormalization group will give the "exact" value as $\frac{1}{4}$ (personal communication). We do not believe this will ever happen.

For the mass term we take the ansatz, as the symmetry breaking which is given by Equation (2.4), which follows on from (2.5) and (2.7) where κ is a constant with dimension of mass. We shall introduce a source for the ϕ field in the form of Equation 2.8. In terms of the quantity $(\langle\varphi\rangle^0 = 250 GeV)$ the electron mass M_e turns out to be a very tiny number $\approx 2 \times 10^{-6}$. Such a number is large only for the top quark if its mass is in excess of 100 GeV. Thus the term which gives rise to this, looks like $f\varphi\bar{t}_L t_R$ where $f \geq \frac{2}{3}$. Some physicists (like Y. Nambu) take this as the defining property of the field φ, i.e. φ is considered as a $t\bar{t}$ composite.

According to Lee and Drell , a proposal is considered according to which the masses of the fermions in the standard model are determined by dynamical symmetry breaking rather than being introduced as arbitrary parameters in the Lagrangian, they are determined self-consistently by the requirement that the proper self-energy vanish the fermion mass shell. They find that in the one-loop approximation it is possible to generate a heavy top quark mass dynamically while the other fermions remain massless [3]. We find that nonzero solutions for m_t do exist and they are always greater than 70 GeV for all values of cutoffs $\Lambda = 5 m_z, 10 m_z, 20 m_z$ and is the only Higgs mass explored in the problem.

One may thus be led to conclude that, while a heavy top quark mass can be generated dynamically, the observed masses of the light fermion cannot be generated in this approach.

Recently, Ruiz-Altaba, Gonzalez and Vargas [11] extended this to a two-loop calculation predicting $m_t = 124$ GeV, $m_H = 234$ GeV and the weak mixing angle $\sin^2 \theta_W \simeq .24$.

Our physical motivation is to remove the fermion masses arbitrary parameters of the Standard Model and treat them as parameters determined by the dynamics in the sub-TeV region. Rather than cancel the divergences, they impose and interpret the Standard Model $\mathcal{L}$ as an effective Lagrangian in order to determine the physical masses. Thus using Sakita's formulation, we get the result.

$$T_c = \frac{\langle\varphi\rangle}{10^3} \exp(-2/g_{eff}\sigma(1 - 4 \sin^2 \theta)) \approx 2.5 \times 10^2\ ^\circ K \qquad (6.3)$$

The exponential factor gives $\exp -26 \approx 10^{-10}$. Assuming that $\sigma(0) \approx m_Z^2$ we obtain $g_{eff}\sigma(0) \approx 1$.

7. Now we go on to the supersymmetry case. The symmetry resolves the dilemma of heirarchy. In a theory possessing fermion–boson symmetry because of the presence of both fermionic and bosonic radiative loops, the radiative shifts for scalar fields are not quadratically divergent. In the limit of broken supersymmetry, the shifts in v^2 only depend logarithmically on the cutoff Λ.

To guarantee that $g_{eff}\sigma(0) \approx 1$, we must take some mass in the theory which is large enough. The simplest mass is obtained by taking supersymmetry such that

$$\phi_0 \rightarrow \phi_0 + \alpha[\Delta M^2]\ \ell n \frac{\Lambda}{\phi_0}\,, \qquad (7.1)$$

where ΔM^2 is the mass splitting between supersymmetry fermions and is of order 1 TeV [12] which gives $\frac{\Delta M^2}{92 m_z^2} \approx 1$ and thus $g_{eff}\sigma(0)$ of the order of unity once again.

This paper started with the possibility of defining a science of life based on atomic physics. What we have shown is that if the model is Z_0 it is possible to have the quark instead of the nucleons. The mass of Z_0 is ≈ 100 GeV, while the mass of the quarks is ≈ 200 GeV so that if Z_0 loops are considered the $t\bar{t}$ quarks go through these loops. The g_{eff} which is proportional to m_Z^{-2} then is of the right magnitude to be able to let these quarks make composites which are manifested through their passage through these loops.

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Table 1

oscillator number	square well	spectral term	spin term	number of states	shells	total number
0	$1s$	$1s$	$1s_{1/2}$	2	2	2
			$1p_{1/2}$	4		
1	1p	$2p$	$1p_{3/2}$	2	6	8
	1d	3d	$1d_{3/2}$	6		
2			$1d_{3/2}$	4	12	
	$2s$	$2s$	$2s_{1/2}$	2		20

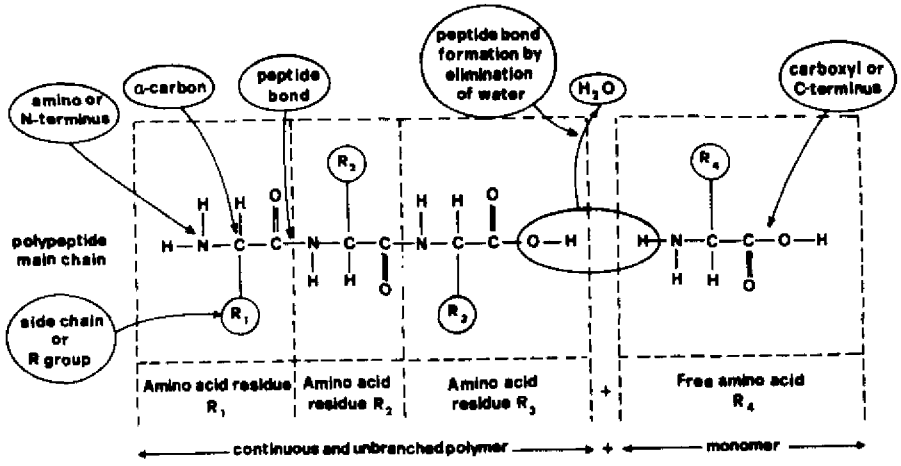

Fig.1

Figure Caption

This diagram is taken from *"From Cells to Atoms - An illustrated introduction to Molecular Biology"*, by Anthony R. Rees and Michael J.E. Sternberg (Blackwell Scientific Publication) p. 13.