J.Phy.Ast.
Print-ISSN: 2320-6756
Online-ISSN: 2320-6764

Journal of
Physics &
Astronomy
WWW.MEHTAPRESS.COM

Full Paper

Stefan Mehedinteanu

CITON - Center of Technology and Engineering for Nuclear Projects, Str. Atomistilor No. 409, BOP-5204-MG-4, Oras Magurele, Ilfov, (ROMANIA)
E-mail: mehedintz@yahoo.com
PACS: 11.15.Ex; 12.38.Mh; 13.38.Be

Received: June 25, 2013
Accepted: August 31, 2013

*Corresponding author's
Name & Address

Stefan Mehedinteanu
CITON - Center of Technology and Engineering for Nuclear Projects, Str. Atomistilor No. 409, BOP-5204-MG-4, Oras Magurele, Ilfov, (ROMANIA)
E-mail: mehedintz@yahoo.com

Numerical analysis of pairs creation by Schwinger effect in nucleons and the beta-decay process acceleration

Abstract

In the paper is done an application of the Dual Ginzburg-Landau-Pitaevski (DGLP) theory to a nucleon, in order to obtain data for the strong electromagnetic fields inside the nucleons and theirs interactions.
Therefore, it is proved for the first time, that in the nucleon exist sufficiently high electromagnetic fields that permit to continue (with a rate of ≅ 1 pair) extract from vacuum of pairs $e^+ - e^-$ (virtual) of high energy electrons, of $W^{\pm}$, Higgs bosons, quarks, by a Schwinger effect, etc, to transform its into real one of very short time life, just like in a veritable laboratory. Thus, it was discovered for the first time that *v.e.v.* is in fact the Schwinger critical field E_{cr} for the pair *W* creation from vacuum. These pairs decay or annihilate into electrons, which passes the monopole condensate barrier as beta-electrons by *quantum tunneling* due of the *phase slip* with $2\pi - \varphi$ and of a Φ_0 energy release, the entire model is proved for a free neutron decay life-time.
Equally, the same Schwinger pairs-production rates are enhanced by a thermal Boltzmann factor in place of quantum tunneling, when this thermalization due of the incidence of an high thermal spike of a photon with nucleons destroys the superconductivity.
This effect is proved in the case of ^{26}Al, through its β-decay to 1.809 MeV γ-ray, when at high temperatures (T_9 = 0.42*GK*) equilibrium is reached between ^{26gs}Al and ^{26m}Al which is relevant to some high temperature astrophysical events such as novae.
In other applications, as based on these data, there are calculated: the Higgs boson energy release due of two gluons fusion during the *pp* collision at LHC, the gluon pair production from space-time dependent chromofield due of the collision of *pp* and of heavy nuclei.

Key Words

Beta decay; Photonuclear reactions; High energy lasers; W,Z,H bosons; G-L theory; Schwinger effect pairs creation; ELI; Gluons; Monopoles condensate.

INTRODUCTION

More, due to the interaction of physics and astrophysics we are witnessing in these years a splendid synthesis of theoretical, experimental and observational results originating from three fundamental physical processes. They were originally proposed by Dirac, by Breit and Wheeler and by Sauter, Heisenberg, Euler and Schwinger.

The vacuum polarization process in strong electromagnetic field, pioneered by Sauter, Heisenberg, Euler and Schwinger, introduced the concept of critical electric field $\mathbf{E_c = m_e^2c^3/e\hbar}$, m_e-electron mass. It has been searched without success for more than forty years by heavy-ion collisions in many of the leading particle accelerators worldwide.

Now, the QCD-monopole has an intrinsic structure relating to a large amount of off-diagonal gluons around its center, similar to the 't Hooft-Polyakov monopole[1,2]. At a large scale where this structure becomes invisible, QCD-monopoles can be regarded as point-like Dirac monopoles.

In the Maximally Abelian (MA) gauge, the off-diagonal gluon contribution can be neglected and monopole condensation occurs at the infrared scale of QCD. Therefore, the QCD vacuum in the MA gauge can be regarded as the dual superconductor described by the DGL theory, and quark confinement can be understood with the dual Meissner effect.

Therefore, in the first part we proceed to a review of our analytical model based on the DGLP theory, already presented in[3], and here we insist more on the equivalence of our model with that described in the works[4-10], from RCNP-Japan, and where, also, is proved the connection between QCD and the dual superconductor scenario.

In the next parts, as based on these data obtained, are calculated: the Higgs boson energy release due of two gluons fusion during the *pp* collision at LHC, the gluon Pairs and quarks pairs production from space-time dependent chromofield; the $e^+ - e^-$ pairs creation due of the thermally-induced vacuum instability as produced by a laser pulse in a the crossed field generated by a single plane wave generated by a single photon, in high energy collisions where jets are the signatures of quark and gluon production.

It will be demonstrated based on the results of DGLP theory, respectively: the value of the maximum chromoelectrical field of the Giant Vortex (*GV*) of $E_0 = 2.18 \times 10^{28}\ N / C$), and of magnetic field of $B = 1.28 \times 10^{17}\ J / Am^2$ due of the monopole condensate current inside the nucleon or of spin-orbit interaction. These values are shown as being near of Schwinger critical electric field and of parallel magnetic field for $e^+ - e^-$ pairs creation by Schwinger effect, all that making possible of *one pair per nucleon* to be obtained. This pair supplies the charges balance (now, not very clear) making possible the quarks conversion ($u \rightarrow d$).

Thus, a new understanding of beta decay process it will be proposed, when, also, a pair of boson $W^- - W^+$ is simultaneously created due of the Schwinger effect in the giant vortex (*GV*) where the electrical field is $E \cong 6.58 \times 10^{29}\ N / C$, and near equally with $E_0 \geq v.e.v. = \mathbf{E}_{cr}^{w^\pm} = 3.5 \times 10^{28} \mathrm{N / C} \leftrightarrow 247 GeV$.

Also, it is given for the first time the demonstration of the discovery, that *v.e.v.* is in fact the Schwinger critical field E_{cr} for a pair of $W^\pm$ creation from vacuum.

This pair decays in beta-electrons during *quantum tunneling* due of the *phase slip* with $2\pi - \varphi$ and of a Φ_0 energy release, and this ad-hoc *bias current* produces a spontaneous suppression of the superconducting order parameter, all the model is proved for a free neutron decay.

Also, it is shown that, equally, the same Schwinger pair-production rate is enhanced by a thermal Boltzmann factor, when the quantum tunneling is substituted by a thermalization which destroy the superconductivity due of the incidence of an high thermal spike of a photon with valence nucleons.

For that, is given a numerical application, when is considered the case of ^{26}Al, through its β-decay to 1.809 MeV γ-ray, when at high temperatures ($T_9 = 0.42 GK$) equilibrium is reached between ^{26gs}Al and ^{26m}Al which is relevant to some high temperature astrophysical events such as novae, this being proved by our model.

THE DGL MODEL FOR NUCLEON SUBSTRUCTURE

In our model is adopted a basic *dual* form of DGDP theory[3] which generalizes the London theory to allow the magnitude of the condensate density to vary in space. As before, the superconducting *order parameter* is a complex function $\boldsymbol{\psi}(\vec{\mathbf{x}})$, where $|\boldsymbol{\psi}(\vec{\mathbf{x}})|^2$ is the condensate density n_s. Also is defined the wave function

$$\boldsymbol{\psi}(\vec{\mathbf{x}}) = \sqrt{\mathbf{n}_s}\ \mathbf{exp(i\boldsymbol{\varphi}(\vec{x}))} \quad (1)$$

or QCD monopole field $\chi_\alpha(\alpha = 1,2,3)$[4,5,11], and where n_s is the London (bulk) condensate density, and φ are real functions describing the spatial variation of the condensate.

The characteristic scale over which the condensate density varies is ξ, the G-L coherence length or the vortex core dimension. The x denote the radial distance of points from the *z*-axis, the superconductor occupying the half space $\mathbf{x} \rhd \mathbf{0}$. Outside of the superconductor in the half space $\mathbf{x} \lhd \mathbf{0}$, one has

$$B = E = H = H_0 \quad (2)$$

where, “the external” vector H_0 is parallel to the surface and correspond to E the external color-electric field inside the hadron flux tube assumed as

$$E = (E_3 T_3 + E_8 T_8) \quad (3)$$

which is formed between valence quarks for the $\mathbf{q\bar{q}}$ pair creation rate[4,5,11]. The Ψ theory of superconductivity[12] is an application of the Landau theory of phase transitions to superconductivity. In this case, some scalar complex ψ function fulfils the role of the order parameter.

First of all, we write the magnetic induction the

$$B = curl A = \nabla \times A \quad (4)$$

where A is the electromagnetic field potential, or the diagonal gluon

$$\overline{\mathbf{A}}_{\boldsymbol{\mu}} = (\mathbf{A}_3^{\boldsymbol{\mu}}, \mathbf{A}_8^{\boldsymbol{\mu}}) \quad (5)$$

and the dual gauge field

$$\overline{\mathbf{B}}_{\boldsymbol{\mu}} \equiv (\mathbf{B}_3^{\boldsymbol{\mu}}, \mathbf{B}_8^{\boldsymbol{\mu}}) \quad (6)$$

as in[4,5,11].

Finally, one gets for the superconducting current density

$$\mathbf{j_s = n_s v_s e = \frac{e\hbar}{2m} n_s \left(\nabla\varphi - \frac{2e}{\hbar} A \right)} \text{ or} \tag{7}$$

$$\mathbf{A = \frac{m}{e^2 n_s} j_s - \frac{\hbar}{2e}\nabla\varphi = \frac{\lambda^2}{2c^2\varepsilon_0} j_s - \frac{\hbar}{2e}\nabla\varphi} \tag{8}$$

$$\mathbf{g = 1/2e = \frac{\alpha}{2} e} \tag{9}$$

Or from[6], it is $\mathbf{B_\mu} = \frac{\mathbf{1}}{\mathbf{2\hat{g}^2}} \frac{\mathbf{k_\mu}}{\phi^2} - \frac{\mathbf{1}}{\mathbf{\hat{g}}} \partial_\mu \mathbf{f} = \frac{\mathbf{1}}{\mathbf{2\lambda_L^{-2}}} \mathbf{k_\mu} - \frac{\mathbf{1}}{\tilde{\mathbf{g}}} \partial_\mu \mathbf{f}$,

that resulting from the following correspondence:

$$B_\mu \to A;\ n_s \to {}^2 \to v^2 \tag{10}$$

$$\mathbf{m_{w\pm}} = \hat{\mathbf{g}}\mathbf{v} = \lambda_L^{-1} \tag{11}$$

$$\partial_\mu f \to \nabla\varphi$$

$$\hat{\mathbf{g}}^2 = \mathbf{4\pi\alpha};\ \mathbf{e^2 = \hat{g}^2 \sin^2\theta_W} \tag{12}$$

For $\mathbf{v = 247GeV \to \lambda_L^2 = \left(4\pi\alpha v^2/(\hbar c)^2\right)^{-1} \cong 5.5e - 36m^2}$, or $\lambda_L \cong 2.3e - 18[m]$, which is the Compton length for $W^\pm$ bosons, see the section 3.2, below. Therefore, a perfect equivalence exists between both models.

The final expression for the free energy then takes the form[3]:

$$f = f_n + \int \left\{ \frac{\hbar^2}{2m} \left| \left(\nabla - i\frac{2e}{\hbar} A \right) \psi \right|^2 + a|\psi|^2 + \frac{b}{2}|\psi|^4 + \frac{B^2}{8\pi} \right\} dV \tag{13}$$

Here, the magnetic induction must be expressed as in (11). One can obtain the basic equations of *DGLP* theory by varying this functional with respect to *A* and ψ*. Carrying first variation with respect to *A*, we find after a simple calculation:

$$\delta f = \int \left[c\frac{ie\hbar}{2m}(\psi^*\nabla\varphi - \psi\nabla\psi^*) + \frac{2e^2}{m}|\psi|^2 A + \frac{curlB}{4\pi} \right] \delta A dV + \int div(\delta A \times B)\frac{dV}{4\pi} = 0 \tag{14}$$

The second integral can be transformed into an integral over remote surface and disappears. To minimize the free energy, the expression in the brackets must be equal to zero. This results in the Maxwell equation

$$curlB = \frac{4\pi}{c} j_s = \frac{1}{c^2\varepsilon_0} j_s (in\_SI) \tag{14.1}$$

or

$$\nabla \times \nabla \times A = \frac{4\pi}{c} . j_s = \frac{1}{c^2\varepsilon_0} j_s \tag{15}$$

provided that the current density is given by

$$j_s = \frac{ie\hbar}{2m}\left(\psi^*\nabla\psi - \psi\nabla\psi^*\right) + \frac{2e^{*2}}{m.}|\psi|^2 .A \tag{15.1}$$

In[6], the equivalent construction being:

$$\partial^\nu {}^*F_{\mu\nu} \equiv k_\mu = -i\hat{g}(\chi^*\partial_\mu\chi - \chi\partial_\mu\chi^*)$$

$$+ 2\hat{g}^{*2} B_\mu \chi^* \chi \tag{23}$$

According to the definition of n_s we can substitute $\psi = \sqrt{n_s}\exp(i\varphi)$. Then (15.1) becomes

$$\mathbf{j_s = \frac{\hbar e}{2m} |\Psi|^2 \left(\nabla\varphi - \frac{2e}{\hbar} A \right)} \tag{16}$$

Equation (16) coincides with (8). This justifies our identification of $|\psi|^2$ with n_s. Variation of (13) with respect ψ* gives, after simple integration by parts,

$$\delta f = \int \left[-\frac{\hbar^2}{4m}\left(\nabla - i\frac{2e}{\hbar} A \right)^2 \psi + a\psi + b|\psi|^2\psi \right] \delta\psi^* dV + \frac{\hbar^2}{4m} \oint \left(\nabla\varphi - i\frac{2e}{\hbar} A\psi \right) \delta\psi^* \cdot dS = 0 \tag{17}$$

The second integral is over the surface of the sample. The volume integral vanishes when

$$\mathbf{-\frac{\hbar^2}{4m}\left(\nabla - i\frac{2e}{\hbar} A \right)^2 \psi + a\psi + b|\psi|^2\psi = 0} \tag{18}$$

Equations (15) and (18) form the complete system of the DGLP theory.

In equation (16), to emphasize: $\lambda = \left(\frac{\varepsilon_0 . m_e . c^2}{n_s . e^2} \right)^{1/2} = 1.17e - 16[m]$, I did a lot of multiplications, and I used the quantized flux: $\Phi_0 = \frac{\pi\hbar}{e}$, and $|\Psi|^2 = n_s$; $n_s = 3$ *_monopoles/V**1.*e* – 45*m*³, *V* = 4/3*πr*³ = 1.45*e* – 45*m*³, *r* = 0.7[*fm*] ε_0 = 8.8*e* – 12[*C*².*N*⁻¹.*m*⁻²].

Since, the magnetic charge of monopole being[13] $g_d = 4\pi\varepsilon_0 \frac{\hbar c}{2e} = \frac{4\pi\varepsilon_0\hbar c}{e^2} e = \frac{137}{2} e = 68.5e$, and assuming that the classical electron radius be equal to "the classical monopole radius" from which one has the monopole mass $m_M = g_d^2 m_e / e^2 = 4700 m_e$, the value of λ *remains unmodified*. In appendix A is presented a fully alternative for calculation of monopole mass.

We now consider the phase transition in superconductors of the second kind.

For this we can omit the non-linear ($|\psi|^2\ \psi$) term in (18), we have

$$\mathbf{-\frac{\hbar^2}{4m}\left(\nabla - i\frac{2e}{\hbar} A \right)^2 \psi + a\psi + b|\psi|^2\psi = 0}$$

$$\mathbf{\frac{1}{4m}\left(-i\hbar - \frac{2e}{1} A \right)^2 \psi = |a|\psi} \tag{19}$$

This equation coincides with the Schrodinger equation for a particle of mass $2m$ and charge $2e$ (in the case of dual, the factor 2 for the charge, which is specific to the "pairs", it is actually 1) in a magnetic field H_0 (in our case the chromo-electrical flux $E(0)$). The quantity $|a|$ plays the role of energy ($E\psi$) of that equation. The minimum energy for a such particle in a uniform electro-magnetic field is

$$\varepsilon_{(0)} = \frac{1}{2}\hbar\omega_B = \frac{1}{2}\hbar\frac{2eH_0}{8mc}[J],$$

H_0-an "external" electro-magnetic field of a dipole created by the pair $u\bar{u}$ (the chromoelectrical colors field)

$$H_0 = E_0 = \frac{de}{4\pi\varepsilon_0 r^3} = 8.33e24\left[\frac{N}{C}\right] \tag{19.1}$$

where is supposed $r \cong 0.05[fm]$-as the electrical flux tube radius, $d = 0.48[fm]$-the distance between the two quarks charges, $m = m_e \cong 9.e-31[Kg]$, usually $H[A/m]$, but here is used as $B = \mu_0 H\left[\frac{J}{Am^2}\right]$

To note, that if we will go in the range of Compton volume for $W^{\pm}$ bosons pair production in vacuum, respectively of radius $r \cong \lambda_{Compton}^{w^{\pm}} = \hbar/m_w c = 2.3e-18[m]$, it results $E_0 \cong 1.75e28[N/C]$, see more in the section 3.2, below.

Hence, equation (19) has a solution only if $|a| \rhd 2\hbar * eH_0/8mc$, when following power-law conformal map is applied for complex number of the r.h.s of (19), or equivalently if the electro-magnetic field is less than an upper critical field, see figures 1a;1b.

$$H_0 \le \frac{4mc|a|}{\hbar e} \le H_{c2} = \frac{\Phi_0}{2\pi\xi^2} = \frac{\pi\hbar c}{2\pi e\xi^2} = 8.33e24\left[\frac{N}{C}\right] \tag{20}$$

and in terms of

$$B_{c2} = H_{c2} = \frac{\pi\hbar c}{2\pi e\xi^2 \cdot c} = 2.5e16\left[\frac{J}{Am^2}\right] \tag{20.1}$$

with $\lambda \rhd \xi$; $\xi = \frac{\lambda}{\kappa} = \frac{0.117}{1.05} = 0.1114[fm]$, or $\kappa \rhd 1/\sqrt{2} \rhd 1 = 1.05$ (of type II-superconductor).

The electrical field is concentrated inside the tube. At large distances from the tube it is shielded by annular superconducting flowing around the tube. This current is analog of the superfluid velocity field surrounding the vortex lines in the superfluid liquid. We can then picture the mixed state as an array of quantized vortex lines. Such vortex lines were predicted by A.A. Abrikosov in 1957. Their existence is crucial for explaining the proprieties of type II superconductors (dual in our case).

The presence of a vortex line in the center of the tube increases the free energy of the superconducting media. The DGLP equations are solved analytically only for $\lambda \rhd\rhd \xi$ (near T_c this means $\kappa \rhd\rhd 1$), *since, in the MA gauge, the charged gluon (M_{ch}) effects become negligible and the system can be described only by the diagonal gluon component at the long distance as* $r >> M_{ch}^{-1} \cong 0.2fm$ [4,5,6] For the short distance as $r \le M_{ch}^{-1} \cong 0.2fm$, the effect of charged gluons appears, and hence all the gluon components have to be considered even in the MA gauge, see appendix A. Thus, when the electrical flux is applied parallel to the superconducting cylinder, the first flux penetrating should be located along the axis of the cylinder.

Substituting j_s from Maxwell equation, we can rewrite (8) as:

$$\frac{1}{\lambda^2}c^2\varepsilon_0\left(\nabla\varphi\frac{2\Phi_0}{2\pi} - A\right) = j_s$$

From Maxwell equation (in *SI*):

$$curlB = \frac{1}{c^2\varepsilon_0}j_s$$

$$\left(\nabla\varphi\frac{2\Phi_0}{2\pi} - A\right) = \frac{\lambda^2}{c^2\varepsilon_0}j_s = \lambda^2 curlB$$

or

$$A + \lambda^2\nabla \times B = 2\Phi_0\nabla\varphi/2\pi \tag{21}$$

The phase φ in presence of vortex line is not a single-valued function of the coordinates. For a vortex line with minimum flux Φ_0, the phase increase by 2π on traversing a closed contour that enclose the line. Thus the integral along such a contour is

$$\oint \nabla\varphi \cdot dl = 2\pi \tag{22}$$

Integrating (21) we find

$$\oint (A + \lambda^2\nabla \times B)\cdot dl = 2\Phi_0 \tag{23}$$

It is not difficult to check that in the range

$$\lambda \rhd\rhd x \rhd\rhd \xi \tag{24}$$

The second term from l.h.s of (23) gives the main contribution. We take the contour of integration in (23) a circle of radius x. For this geometry the vector ($\nabla \times B$) has only one component $(\nabla \times B)_\varphi$ along the contour.

The integration is then simple and we have

$$(\nabla \times B)_\varphi = -\frac{dB}{dx} = \frac{2\Phi_0}{2\pi x\lambda^2} \tag{25}$$

To note (in *cgs*):

$$(\nabla \times B)_\varphi = \frac{4\pi}{c}en_s v_{s\varphi}$$

Equation (25) then gives $v_{s\varphi} = \hbar/2mx$ for the superfluid velocity as it must be for a vortex line in a superfluid of particles with mass $2m$.

Integrating of (25) for B gives

$$\mathbf{B(x)} = \frac{2\mathbf{\Phi_0}}{2\pi\lambda^2}\log\left(\frac{\lambda}{\mathbf{x}}\right) \tag{26}$$

This equation is valid in the interval (24) with logarithmic accuracy.

Notice also that every vortex carries the flux Φ_0 and hence the mean value of B over the cross-section of the cylinder is

$$\overline{\mathbf{B}} = 2\nu\mathbf{\Phi_0} \tag{27}$$

where ν is the number of lines per unit area. This result is invalid near the upper critical flux H_{c2} where the cores of the vortex lines begin to overlap. To calculate this number we have to take into account the interaction between vortex lines. As the first step we have to find the electrical field trough a loop of arbitrary radius surrounding the line without the restriction (24). Let us calculate the curl of the both sides of (21).

Note that

$$curl\nabla\varphi = 2\pi \cdot n_z \cdot (x) \tag{28}$$

and $curlA = B$, where $\delta(x)$-the Dirac function

Where r is the two-dimensional radius-vector in the $x - y$ plane and n_z is a unit vector along axis z (We assume that the axis of the vortex line coincides with z). Indeed, integrating ∇φ along the contour encircling the line and transforming the integral by Stokes' theorem into an integral over a surface spanning the contour, we have according to (22)

$$\oint \nabla\varphi \cdot \mathbf{dl} = \int \mathbf{curl}\nabla\varphi \cdot \mathbf{dS} = 2\pi \tag{29}$$

Since this equation must be satisfied for any such contour of integration, we have (28).

Finally, we obtain

$$B + \lambda^2 curlcurlB = 2n_z\Phi_0\delta(x) \tag{30}$$

Using the vector identity $curlcurlB = \nabla divB - \Delta B = -\Delta B$, we obtain

$$B - \lambda^2 \cdot B = 2\Phi_0\delta(x) \tag{31}$$

This equation is valid only at all distances

$$\mathbf{x} \rhd\rhd \xi \tag{32}$$

Throughout all the space except the line $x = 0$ equation (31) coincides with the London equation (12)

The $\delta(x)$ function on r.h.s defines the character of the solution at $x \to 0$. Actualy this singularity has already been defined in (26), which is valid at small x.

The solution of this equation at $x \to \infty$ is $B(r) = const \cdot K_0(x/\lambda)$, where K_0 is the Hankel function of imaginary argument. The coefficient must be defined by matching with the solution of (26). Using the asymptotic formula $K_0(x) \approx \log(2/\gamma x)$ for $\mathbf{x} \lhd\lhd \mathbf{1}$, where $\gamma = e^C \cong 1.78$ (C is Euler's constant), we finally have

$$\mathbf{B(x)} = \frac{2\mathbf{\Phi_0}}{2\pi\lambda^2}\mathbf{K_0(x/\lambda)} \tag{33}$$

Exactly, the same solution is obtained by G.Bali et al[12], respectively the equation (2.6).

Using equation (33) we can rewrite (26) as:

$$\mathbf{B(x)} = \frac{2\mathbf{\Phi_0}}{2\pi\lambda^2}\log\frac{2\lambda}{\gamma\mathbf{x}}, \quad \mathbf{x} \lhd\lhd \lambda \tag{34}$$

In opposite limit of large distances one can use the asymptotic expression $K_0(\mathbf{x}) \approx (\pi/2\mathbf{x})^{1/2}\mathbf{e}^{-\mathbf{x}}$ for $\mathbf{x} \rhd\rhd \mathbf{1}$. Thus, at large distances from the axis of the vortex line the field decreases according to

$$\mathbf{B(x)} = \frac{2\Phi_0}{(8\pi \mathbf{x}\lambda^3)^{1/2}}\mathbf{e}^{-\mathbf{x}/\lambda}, \quad \mathbf{x} \rhd\rhd \lambda \tag{35}$$

Accordingly the superconductive current density decreases (in *SI*):

$$\mathbf{j}_\varphi = -\frac{\mathbf{c}}{4\pi}\frac{\mathbf{dB}}{\mathbf{dx}}(4\pi \mathbf{c}\varepsilon_0) = \frac{2\mathbf{c}^2\varepsilon_0\mathbf{\Phi_0}}{8(2\pi^3\mathbf{x}\lambda^5)^{1/2}}\mathbf{e}^{-\mathbf{x}/\lambda}$$

We can now calculate the energy ε of the vortex line. The magnetic part of free energy corresponding to London equation is given by the integral.

$$\mathbf{F_B} = \frac{1}{8\pi}\int\left[\mathbf{B}^2 + \lambda^2(\mathbf{curlB})^2\right]\mathbf{dV}$$

Indeed, by varying the expression with respect to B, we immediately obtain the London equation (12). The main contribution to the integral is due to the second term, which contains a logarithmic divergence. Substituting (25) in the above eq., and integrating in the range (24), we obtain for the energy per unit length of vortex line.

$$\varepsilon = \left(\frac{2\mathbf{\Phi_0}}{4\pi\lambda}\right)^2\log\left(\frac{\lambda}{\xi}\right) \tag{35.1}$$

This equation, explains why only vortex lines with the minimum flux Φ_0 are the most favorable. The energy of a line is proportional to the square of its magnetic flux. Thus, the fragmentation of one line with the flux $n\Phi_0$ into n lines with flux Φ_0 results in an n-fold gain in energy.

The vortex energy to create a W-boson by a Schwinger effect is given as (see section 3.2.):

$$\varepsilon_{vortex} = Vc^2\varepsilon_0 H_{c2}^2/8\pi = 1.16e-08[J] \tag{36}$$

where V-is the volume of one vortex ($1/3V$)

The interaction between vortexes at distance x and of separation $d = x - \lambda$ is given as in[3], see figure 1a;1b.

$$\varepsilon_{\mathrm{int-pair}} = \mathbf{c}^2\varepsilon_0\frac{4\mathbf{\Phi}_0^2}{2^{7/2}\pi^{3/2}\lambda^2}\left(\frac{\lambda}{\mathbf{d}}\right)^{1/2}\mathbf{e}^{-\mathbf{x}/\lambda} =$$

$$\frac{1.e17*8.82e-12*4*2.06e-15^2}{2^{3.5}3.14^{1.5}0.117e-15*0.117}\left(\frac{\lambda}{0.36*\lambda}\right)^{1/2}\mathbf{EXP}\left(-\frac{1.36*\lambda}{\lambda}\right) =$$

$$6.65e-09[\mathbf{J/fm}] \Rightarrow 41.6[\mathbf{GeV}]$$

The energy of the neutral *boson Z* is assimilated with the vortex-vortex two pairs of quarks spins (1/2 + 1/2 = 1) interaction energy $\varepsilon_Z = 2 * \varepsilon_{\mathrm{int}\text{-}pair} = 84[GeV]$, with above $\varepsilon_{\mathrm{int}\text{-}pair}$. When the three pairs of vortexes with the outermost ($d \cong 0.04\,fm$) vortices lines which interacting (repel) at the center of the triangle, that will generate a neutral current in this zone called *Higgs boson* (H). Its spin and charge are both zero due of the vortices coalescence, here into a giant vortex (GV), all that happens during the triangular arrangement of the

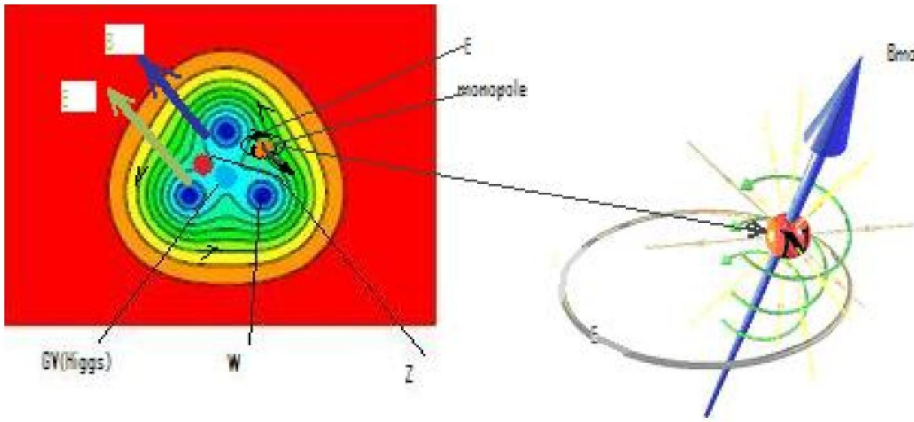

Figure 1a : The Giant-Vortex (after an idea from Ref.[15]) that could be also the arrangement for the nucleon (only illustration).. A spin-orbit nonabelian field is shown (after an idea of the ref[14]).

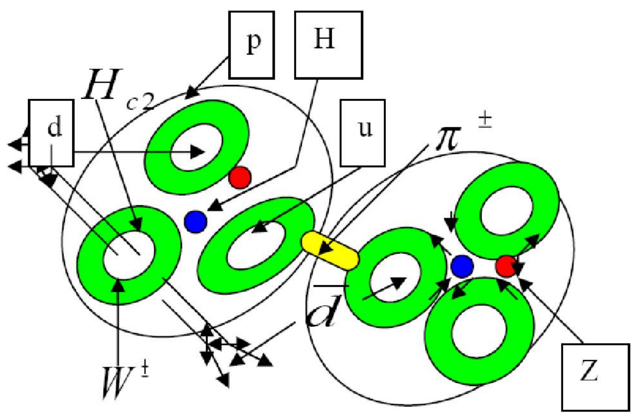

Figure 1b : Abrikosov's triangular lattice for a nucleon (proposal after an idea from[15])

lattice, see figures 1a;1b. Thus, it results another energy state-maximum possible ($d \cong 0$), probable that of *Higgs boson (H)*: $\varepsilon_H = 3 * \varepsilon_{\text{int-}pai} = 3 \times 41.6 \rightarrow 125[GeV]$.

In other words, in order to equilibrate these energies are necessarly to admit the existence of two particles Z, H of exactly these energies values.

THE PAIR CREATION INSIDE THE NUCLEON BY SCHWINGER EFFECT

Electron-positron pair

An interesting aspect of virtual particles (in vacuum) both theoretically and experimentally is the possibility that they can become real by the effect of external fields. In this case, real particles are excited out of the vacuum. In the framework of quantum mechanics by Klein, Sauter, Euler and Heisenberg who studied the behavior of the Dirac vacuum in a strong external electric field. If the field is sufficiently strong, the energy of the vacuum can be lowered by creating an electron-positron pair. This makes the vacuum unstable.

A particle with charge q and mass m in a constant magnetic field undergoes circular motions with the Larmor frequency $\omega_c = qB/m$. In quantum theory charged particles occupy the Landau levels with energy $\mathbf{E} = \hbar\boldsymbol{\omega}_c \cdot \mathbf{n}$ and in a strong magnetic field of $B_c \cong 10^{16}[T]$, the energy difference of Landau levels $\boldsymbol{\delta}\mathbf{E} = \hbar\boldsymbol{\omega}_c$ can be comparable to the rest mass of the particle. The number of Landau levels is $\mathbf{N} \leq \frac{\mathbf{m}\boldsymbol{\omega}_c\mathbf{L}^2}{2\pi\hbar} = \frac{\mathbf{qBL}^2}{2\pi\hbar}$, where L is the length. In the transverse direction of a magnetic field above the critical strength electrons of an atom are strongly bounded and fill the lowest Landau levels but in the parallel direction are attracted by the Coulomb force. As classically the charged particle moves along a spiral of circular motion in the transverse direction and linear.

In the case analyzed in[16,17], in which $E^2 - B^2$ and $E \cdot B$ are not both zero, one can go to a frame in which E and B are parallel with magnitude E and B, and is obtained the imaginary part of the one-loop effective action per four-volume for spinor *QED*.

In a pure magnetic field $B = Be_z$ along the z-direction, $A = (0,Bx,0)$, and for the electric and magnetic fields parallel to each other along the z-direction, the 4-potential is given by $A_\mu = (-Ez,0,Bx,0)$, that copy very well with our new understanding, when the electric field E is that of the quark-anti-quark pairs, and B is induced either by the spin-orbit interaction of the monopole (nonabelian field), or by the monopole current as Rashba effect, see appendix A, and figure 1a.

In[16] is obtained the pair-production rate for fermions (eq. (67)) as

$$\frac{\Gamma_{JWKB}}{V}=\frac{\alpha}{\pi\hbar}\left(\sqrt{4\pi\varepsilon_0}\right)^2(Bc)\cdot E\cdot\coth\left(\frac{\pi Bc}{E}\right)\exp\left(-\frac{\pi E_c}{E}\right)\quad(73)$$

for spin –1/2 particle,

where $\alpha = 1/137$, $V[m^3 s]$ and $\sqrt{4\pi\varepsilon_0}$ to convert $cgs \to SI$ (to note that the authors of these articles do not have used this factor, that affecting *enormous* the numerical results with $\cong 10^{-11}$), JWKB meaning (Jeffreys-Wentzel-Kramers-Brillouin) model[16].

The discrete spectrum due to the magnetic field is the Landau levels for charged particles. Note that all the Landau levels are non-degenerate for the scalar particles, whereas all the states of the Dirac spinor are doubly degenerate, j, $\sigma_+ = 1/2$, and $j - 1$, $\sigma_- = -1/2$, except for the unique lowest Landau level, $j = 0$, $\sigma_+ = 1/2$.

The effect of magnetic field is the same as shifting the effective mass $m_*^2c^4 = m_e^2c^4 + qB\hbar c^2(2j+1-\sigma)$ for fermions for each Landau level.

(a) Case 1 - The use of a nonabelian field

The Compton space-time volume of an electron has the size

$V_{Compton} = \lambda_C^3 \times (\lambda_C/c) = 1.59e-70m^3s$,

Where $\lambda_C = \hbar/m_*c = 4.6e-16[m]$, the effective mass is $m_* = \sqrt{m_e^2c^4 + qB\hbar c^2}/c^2$, the *critical field* $E_c = \frac{m_*^2c^3}{e\hbar} \cong 8.58e23 < E = 8.24e24[N/C]$; $B = 2.86e15[T]$ as from eq.(A.40), that results $m_* = 7.14e-28kg$

The electron energy is $\varepsilon = m_*c^2 \cong 0.4GeV$- the same and sufficiently to broken the quark-anti-quark string strength $\sigma \cong 0.4GeV$ during beta-decay process followed by the release of a beta-decay electron from $W^{\pm}$ decay, and which gets the final beta energy as equally to that of the out of barrier turning point after the tunneling and accounting for the valence nucleons interactions (shell-energy levels). The number of *assaults* of the barrier, like in Gamow theory[18] is $n_a = v_b/R$; where the velocity is $v_b \cong (2\varepsilon/m_*)^{1/2}$, and the radius of the barrier is $R \cong \lambda_C$; $\varepsilon = \hbar eB/m^* \cong 0.4GeV$ the energy of the particle for the first Landau level (as above), and we can see that it results to be equally with the rest mass of the electron, that resulting $n_a \cong 9.e23s^{-1}$. In case of *WKB*[18], the transmission coefficient $T = 2\frac{\sqrt{2m|V-Q|}}{\hbar}\Delta r$, and the decay constant $R = n_a e^{-T}$.

For the thick barrier the transmission coefficient is $T = 2\pi\frac{Qb}{\hbar v} = 2\pi\frac{\sqrt{2mQ}}{\hbar}b$; where, the kinetic energy of the particle after the barrier at b is $Q \equiv \frac{1}{2}mv^2$

To "materialize" a virtual $e^+ - e^-$ pair in a constant electric field E the separation d must be sufficiently large $eEd = 2mc^2$

Probability for separation d as quantum fluctuation

$$P \propto \exp\left(-\frac{d}{\lambda_{Compton}}\right) = \exp\left(-\frac{2m^2c^3}{e\hbar E}\right) = \exp\left(-\frac{2E_{cr}}{E}\right)$$

The emission (transmission through barrier) sufficient for observation when $E \approx E_{cr}$, with $Q = eE_{cr}d$, results

$$T = 2\pi\frac{mcb}{\hbar} \cong \frac{2\pi b}{\lambda_C}, \text{ or } b \cong \lambda_C/2\pi \quad (73.1)$$

With these values it result : the number of pairs $\Gamma_{JWKB}/s = 4.65e23s^{-1} \approx n_a$, for a volume $V \cong (\lambda_b)^3 \cong 1.0e-46[m^3] \geq V_b$, the penetration length $\lambda = 0.117\,fm$, and for a four-volume of $\lambda_b^4/c \cong 1.6e-70[m^3s]$, results $\Gamma_{JWKB} \cong 0.72 \approx 1pair$, or in other words, all the time inside the nucleon is *available* one real pair of electron-positron which combine with quarks pairs, as was shown before, resulting an e^+, or e^- which help the quarks transformation ($u \to d$) for beta-decay. In our model, $W^{\pm}$ is also created (see bellow) there as a *vortex* which decay into an electron which takes the energy at the *turning point* out of the barrier equally with the binding energy of nucleon in isotope nucleus, and it *passes* the barrier of monopole condensate characterized by an *quantum tunneling* suppression given as: $\exp(-\Delta E\tau/\hbar) \cong 3.96e-27$, where, as $\tau \cong \hbar/m_*c^2 \cong 1.56e-24[s]$ near $\tau_{GL} = \pi\hbar/(8k_BT_c) = 1.5e-24[s]$ is the Ginzburg-Landau life time of $W^{\pm}$ bosons, since these decay in beta electrons, in the same time with pair generation, their lifetime need to be near equally with the *lifetime* of the pair $\tau_{e^+-e^-} = \hbar/m_*c^2 \cong 1.56e-24[s]$, that is happen, and ΔE, which corresponds to the height of monopole condensate barrier, due of the *phase slip* with $2\pi - \varphi$ and of a Φ_0 energy release: $\Delta E = c^2\Phi_0^2\varepsilon_0/d_b$; $d_b \cong 7.15\lambda$, and $\Delta E = 3.898e-09[J]$; and the quantized flux is :

$\Phi_0 = \pi\hbar c/e \to \text{usually}\ \frac{\pi\hbar}{e} = 2.07e-15[Tm^2]$. To note that ΔE

Thus, the probability (rate), into a more simple way- without the external interactions of the neutron (free-not bounded), is s given as:

$\Gamma_{JWKB}V\exp(-\Delta E\tau/\hbar) \cong 1.84e-03s^{-1} \to \tau_{1/2} \cong 544[s]$, that corresponds for *free neutrons decay* by emission of an electron and an electron antineutrino to become a proton, $\mathbf{n^0 \rightarrow p^+ + e^- + \bar{v}_e}$, with half-life of 611*s*.

(b) Case 2 - The use of Rashba effect

The Compton space-time volume of an *heavy* electron (gluon!) has the size

$\mathbf{V_{Compton} = \lambda_b^3 \times (\lambda_b/c) = 8.e-74m^3s}$,

Where $\mathbf{\lambda_b = \hbar/m_* c = 7.e-17[m]}$, the effective mass is

$$\mathbf{m_* = \sqrt{m_e^2c^4 + qB\hbar c^2}\,/c^2}\ ,\ \mathbf{m^* = 4.77e-27[kg]} \qquad (38)$$

and the *critical field*

$$\mathbf{E_c = \frac{m_*^2c^3}{e\hbar} \cong 3.84e25 < E = 2.18e28[N/C]}$$

The "heavy electron" energy is $\mathbf{\varepsilon = m_* c^2 \cong 2.68GeV}$ - the same and sufficiently to either broken or to create an quark-anti-quark string, that fact it will be established experimentally!

Therefore, if we consider the situation of the *GV*, in eq. (37) is introduced the magnetic field induced by the entire monopole condensate as being $\mathbf{3\times\varepsilon_{vortex} = \frac{V\varepsilon_0 c^2 B^2_{monop}}{8\pi} \rightarrow 7.9e-07[J]}$, where $\mathbf{B_{monop} \cong E_{monopol}/c = 1.28e17[J/Am^2]}$;

$$\mathbf{E^2_{monop}} = \frac{(\varepsilon_{monop})^2}{\varepsilon_0(\lambda_C^{e\bullet})^2\hbar c} \qquad \mathbf{(38.1)}$$

$\mathbf{\varepsilon_{monop} \cong H_R\_or\_ \cong 3monopoles\_mass}$, the Rashba energy being $\mathbf{H_R = 1.2e-09[J]}$ from eq. (A.32), and $\mathbf{3\times\varepsilon_{vortex}}$ from (36), resulting $\mathbf{E_{monopole} \cong E_{cr}^{e^+-e^-} \cong 3.835e25[N/C]}$, and the electrical field of the Giant Vortex (*GV*), (see figures 1a;1b) is $\mathbf{E^2 = \frac{(3\times\varepsilon_{vortex})^2}{\varepsilon_0(\lambda_C^{e*})^2\hbar c} \rightarrow E = 2.18e28[N/C]}$, where $\mathbf{V \cong 1.45e-45m^3}$ for the Giant Vortex, near the value of $\mathbf{v.e.v = 3.5e28[N/C]}$. To mention that, independently we can obtain the critical value of magnetic field as: $\mathbf{B^{cr}_{monop} \cong E^{e^\pm}_{cr}/c \cong \frac{3.84e25}{2.997e8} \equiv 1.28e17[J/Am^2]}$

With these values it results: the number of pairs created $\mathbf{\Gamma_{JWKB}/s = 4.12e24[s^{-1}] \approx n_a}$, and finally with Compton volume it results $\mathbf{\Gamma_{JWKB} \cong 1pair}$, or in other words, all the time inside the nucleon is *available* one real pair of electron-positron which combine with quarks pairs as was described above, or creates a new quarks pair-for a new hadron type (with four quarks as a new state of mater!).

In our model, in the same time is created a pair $W^\pm$ by a Schwinger effect (see the next section).

This "heavy electron" itself can not *pass* the barrier of monopole condensate as characterized by an *quantum tunneling* suppression given as: $\mathbf{\exp(-\Delta E\,\tau/\hbar) \cong 1.2e-04}$, where, as $\mathbf{\tau \cong \hbar/m_* c^2 \cong 2.33e-25[s]}$ small that $\mathbf{\tau_{e^\pm} = 1.5e-24[s]}$, or the Compton length being too small by the barrier width $\mathbf{\lambda_C \cong 7.e-17[m] << b \cong 7\lambda = 0.83fm}$. The same conclusion is obtained if we apply eq. (37.1). Therefore, this passes only when a external energy (by collisions!) is applied in order to maintain "open" the barrier by "melting" it.

Therefore, it look like of two permanent pairs of $e^+ - e^-$ and of $W^\pm$ to be present during the beta decay, or an energy of a vortex $W^\pm$ exists here, as it was found in the previously work[3]. In the same way, many others particles could be created here, like *s*; *c*; *t* quarks ; muons etc., but anyone can not penetrates the barrier as itself, these can only decay, or annihilate into electrons of low energy and neutrinos.

To remember that in[3], we have described the process as a spontaneous nucleation of a normal-state belt across the strip with 2π-φ *phase slip* with Φ_0 release, and a *bias currents* which may produce a spontaneous suppression of the superconducting order parameter (ψ), and when a vortex ($W^\pm$) (after it is created by the Schwinger effect, as the main finding of this new work, see section3.2.), crossing from one strip edge (just the monopole condensate barrier) to the opposite one induces a *phase slip* without creating a normal (vacuum) region across the strip (one of three vortexes of nucleus) width. Then, is treating the vortex as a particle moving in the energy potential formed by the superconducting currents around vortex center inside the strip and by the Lorentz force induced by the bias current. For a free neutron the *dark count rate* in[3] is obtained as $\mathbf{R_{dark}} = \mathbf{7e-04s^{-1}}$, which is comparable with the above result.

To see the order of magnitude we extract from[3] some results on the decay of a free neutron, thus, the bias current is:

$$I = \frac{2w}{\pi\xi} I_0 \kappa(1-\kappa^2)$$

and the main superconducting current is:

$$\mathbf{I_0 = \varepsilon_0 \frac{c\Phi_0}{8\pi\Lambda}; \Lambda = \frac{2\lambda^2}{h_z}}$$

Here, $h_z \cong \lambda$-the axial (z) height of the monopole condensate.

Also, here, the critical current at which the energy barrier vanishes for a single vortex crossing:

$\mathbf{I_c = \frac{2\mu^2 w I_0}{2.72\pi\xi}}$; respectively:

$I = 5.197e3$; $I_c = 1.6e4$; $I_0 = 3.e5$; where $\mu^2 = 1 - \kappa^2$; $\mathbf{w/\xi = 4.15}$; $\mathbf{\kappa = \lambda/\xi \cong 0.117/0.1114}$; and $\mathbf{I/I_0 \cong 0.017 \to 1e/68e \cong 0.014}$

The Bosons pair production

(a) Case 1 - W± bosons creation

It has long been known that an inevitable consequence of Dirac's theory of the electron is that in regions of sufficiently high energy density, the quantum vacuum can break down in a spontaneous generation of electron-positron pairs. Following the initial results of Sauter, Heisenberg and Euler and Weisskopf, in a seminal work, Schwinger derived a central result of strong-field quantum electrodynamics, the rate per unit volume of pair creation R in a constant and uniform electric field of strength E, of leading order behavior,

$$R = (E/E_{cr})^2 (c/\lambda^4)(8\pi^3)^{-1} * \exp(-\pi E_{cr}/E)$$

for $E/E_{cr} << 1$, positron charge e, mass m, Compton wave-length $\lambda = \hbar/mc$ and so-called "critical" electric field $E_{cr} = m^2c^3/e\hbar$.

Now, the energy corresponding to E_{cr} is given as: $\mathbf{v^2 = E_{cr}^2 \varepsilon_0 \left(\lambda^{W^\pm}_{Compton}\right)^2 \hbar c}$, or $\mathbf{v^2 = \frac{m_W^4 c^6}{e^2\hbar^2}\frac{\hbar^2}{m_W^2 c^2}\frac{4\pi\hbar c \varepsilon_0}{4\pi} = \frac{m_W^2 c^4 \varepsilon_0 \hbar c 4\pi}{4\pi e^2} = \frac{M_W^2}{4\pi\alpha} \cong \frac{0.25 M_W^2}{\pi\alpha} \cong (267 GeV)^2}$, which in fact is the vacuum expectation value (*v.e.v.*), here the fine-structure constant is $\mathbf{\alpha = \frac{e^2}{4\pi\varepsilon_0 \hbar c} = \frac{1}{137}}$

If we will remember the Fermi disintegration constant

$$\mathbf{G_F = \frac{1}{4\sqrt{2}}\frac{g^2}{M_W^2} = 1.15 \times 10^{-5} GeV^{-2}}$$

where the mass of boson $\mathbf{M_W = m_W c^2 = 81 GeV}$, and $\mathbf{e = g \sin\theta_W}$; the Weinberg angle $\mathbf{\sin^2\theta_W \cong 1/4}$, and the Higgs vacuum expectation value *v.e.v* or $\mathbf{v = \left(\sqrt{2} G_F\right)^{-1/2} = 247 GeV}$, if we substitute $\mathbf{M_W = \frac{gv}{2}}$ from Standard Model.

$$\mathbf{M_W^2 = \frac{e^2}{4 s_W^2} v^2 = \frac{\pi\alpha}{s_W^2} v^2 = \left(\frac{37.2}{s_W} GeV\right)^2 \cong (80 GeV)^2}$$, where

$\mathbf{s_W^2 = \sin^2\theta_W \approx 0.22}$

If in eq. above in place of $0.25 \to 0.22$, it will result $v = 247 GeV$

Therefore, it is obtained for the *first time* the derivation of the *Higgs vacuum expectation value field* (*v.e.v*) as being *Schwinger critical field* for creation of pair of bosons W^- – W^+ as needed in beta-decay process. To note that the Standard Model is centered, also in $W^\pm$, the difference being that, in my model this pair is created by the Schwinger effect, and the field of the nucleon substructure having values which support this effect.

Numerically, $\mathbf{m_W = 81 GeV \to 1.442e-25 kg}$, $\mathbf{r \cong \lambda^{W^\pm}_{Compton} = \hbar/m_W c = 2.31e-18[m]}$ for (GV), results

$$\mathbf{E^2 = \frac{(3 \times \varepsilon_{vortex})^2}{\varepsilon_0 (\lambda_C^W)^2 \hbar c} \to E = 6.58e29[N/C]}$$, and

$$\mathbf{3 \times \varepsilon_{vortex} = \frac{V\varepsilon_0 c^2 B^2}{8\pi} \to 7.9e-07[J] \to \cong 5 TeV}$$, and

$\mathbf{B \cong E_{monopol}/c = 1.28e17[J/Am^2]}$ with E_{monop} from eq. (38.1), and Schwinger critical field $\mathbf{E_{cr} = m_W^2 c^3/e\hbar = 3.5e+28[N/C] \leftrightarrow v.e.v = 267 GeV}$;

The rate of W^- – W^+ pair production is again $R \cong 1$, with the Compton volume $\mathbf{V_C^W = (\lambda_C^W)^4/c = 9.54e-80[m^3 s^{-1}]}$, and $\mathbf{R/s = 1.56e26 s^{-1}}$. Again, due of very small Compton length by comparison with the barrier width, $W^\pm$ itself can not penetrate the barrier, as describe above.

To note, that in order to obtain the pair rate of ≈ *one* for $W^\pm$, only then when are used the values of GV energy $\mathbf{(3 \times e_{vortex} \cong 7.9e-07[J] \to 5 TeV)}$, which can means that these particles are obtained by the *melting* of two gluons (*gg*), as is find in[3], and at LHC-see section 5. Therefore, a collision of a such value can maintain *"open"* the barrier till the release of decay products!.

For the first time in[19] is studied the entire dynamics of energy conversion from initial overcritical electric field, ending up with thermalized electron-positron-photon plasma.

Such conversion occurs in a complicated sequence of processes starting with Schwinger pair production which is followed by oscillations of created pairs due to back-reaction on initial electric field, then production of photons due to annihilation of pairs and finally isotropization of created electron-positron-photon plasma. Evolution of electric field E and pairs bulk parallel momentum for $\mathbf{E > 30 E_{cr}}$, shows that following oscillations E tend asymptotically to E_{cr} after $\mathbf{\cong 1000\tau_c}$. After some time, the photons energy density becomes equal and then overcomes the pairs energy density. This growth continues until the equilibrium between pairs annihilation and creation processes is established $e^- - e^+ \leftrightarrow \gamma\gamma$.

A new understanding of beta decay

In the classic understanding of β^- disintegration $n \to p + e^- + \bar{\nu}_e$, or this occurs when one of the down quarks in the neutron (*udd*) decays into an up quark by emitting a virtual W^- boson, transforming the neutron into a proton (*uud*). The W^- boson then decays into an electron and an electron antineutrino: $udd \to uud + e^- + \nu_e$.

In our *new understanding* of beta-decay process, it results that, in order to make possible this transformation (balance of charges) it needs to have supplementary an e^- + e^+ pair as created by a Schwinger mechanism, as will we show above, and which is necessarily to *always* exist here, and consequently, for the beta decay process, we have in terms of quarks for the neutral mesons: $u\bar{u}; d\bar{d}$:

$$d(-1/3e) + e^+(+3/3e) = u(+2/3e) \qquad (75)$$

$$\bar{d}(1/3e) + e^-(-3/3e) = \bar{u}(-2/3e)$$

and for β^+ decay can only happen inside nuclei when the daughter nucleus has a greater *binding energy* (and therefore a lower total energy) than the mother nucleus. The difference between these energies goes into the reaction of converting a proton into a neutron, a positron and a neutrino and into the kinetic energy of these particles.

In an opposite process to the above negative beta decay, the weak interaction converts a proton into a neutron by converting an up quark into a down quark by having it emit a W^+ or absorb a W^-.

β^+ decay of nuclei (only bounded proton) when: $p \to n + e^+ + \nu_e$, or $\text{energy} + uud \to udd + e^+ + \nu_e$

or,

$$u(2/3e) + e^-(-3/3e) = d(-1/3e) \qquad (76)$$

$$\bar{u}(-2/3e) + e^+(3/3e) = \bar{d}(1/3e)$$

In the process of electron capture, one of the orbital electrons, usually from the *K* or *L* electron shell, is captured by a proton in the nucleus, forming a neutron and an electron neutrino.

$$p + e^- \to n + \nu_e$$

In our *new understanding*, when the vortex equilibrium is disturbed by the transformation of one quark ($d \to u$) due of the interaction with the new created e^- + e^+ pair (coincidently of the same energy as that of the strings (neutral mesons) $u\bar{u}; d\bar{d}$), then, this is accompanied by a creation of a $W^\pm$ boson which decay into $e^\pm$ and neutrino which pass the barrier as *crossing vortex*, that being the new idea of this model.

SCHWINGER PAIR-PRODUCTION THERMALLY STIMULATED BY A LASER PULSE

From[20] results that the Schwinger pair-production rate by a time-dependent electric field is enhanced by a thermal factor of the initial Bose-Einstein distribution known as the Boltzmann factor.

Thus, it is found that the Schwinger pair production rate at finite temperature is enhanced by the thermal Boltzmann factor: $f_{nk}(T) = \frac{1}{e^{\omega_{nk}^{(-)}/T} - 1} \cong \frac{k_B T}{\omega_c \hbar}$.

Where the cyclotron frequency induced by the monopole *B* is ω_c.

Thus, with B = 2.8*e*15[*T*] from eq.(A.40), T = 0.4*e*9*K* for decay of $^{26}Al \to {}^{26}Mg$, results $\omega_c = eB/m_* = 6.4e23 s^{-1}$, finally the enhancement factor $f_{Boltzmann} \cong 8.6e-05$ or the lifetime is reduced with 1.*e*7, and the number of pairs per a thermal spike it could be $\Gamma_{JWKB} \cdot V \cdot \tau \cdot f_{Boltzmann} \cong 1 pair$, if the volume affected by laser is $V\tau \cong (\lambda_C^{e*})^3 * 30 zs[m^3 s]$; represent the turning point at which the decay energy $\cong 1.8 MeV$ for ^{26g}Al.

To note that the duration of the *single attosecond spikes* in the APT amounts to 300 *as*/2 means 30*zs* in the projectile frame[21-23]. This value approaches the natural QED time scale of $1/m \approx 1 zs = 10^{-21}[s]$. Near the same result is obtained in case 2, or when a barrier is "open" by the thermal spike, the heavy electrons passing more easy.

An important thing to be verified experimentally is that by using a thermal spike of only T = 0.4e9K, when it could be obtained one $W^\pm$ pair, and one of H boson.

We may thus conclude that the Schwinger pair-production rate is indeed enhanced by the thermal effect given by the Bose-Einstein distribution expressed by the Boltzmann factor.

In this case, in place of the *quantum tunneling* suppression factor we introduce the *thermal stimulation* with Boltzman factor. The model it needs to be experimentally tested, for example on radioisotope $^{26}Al \to {}^{26}Mg$, as it was described bellow as the proposal to test the model, also suggested in[24-26].

To mention that from the model based on *bias current* (described above) it results a necessary photon flux of $R_h \cong 5.18e13 photons/s$, and from this new model results a duration $\tau \cong 30 zs$; both parameters can be obtained only by the laser from ELI (Extreme Light Infrastructure)[27,28].

THE GLUON PAIR PRODUCTION FROM ARBITRARY TIME DEPENDENT CHROMO-ELECTRIC FIELD VIA SCHWINGER MECHANISM

The subject of quark/anti-quark and gluon pair production from the non-abelian field is relatively new and is not fully solved[29]. It might be important for the production of the quark-gluon plasma (QGP) in the laboratory by high energy heavy-ion collisions. Lattice QCD predicts the existence of such a state of matter at high temperatures (~ 200 MeV) and densities. In high energy heavy-ion collisions at RHIC and LHC[27] the receding nuclei might produce a strong chromofield which would then polarize the QCD vacuum and produce quark/anti-quark pairs and gluons. These produced quarks and gluons collide with each other to form a thermalized quark-gluon plasma. The space-time evolution of the quarkgluon plasma in the presence of a background chromofield is studied by solving relativistic non-abelian transport equation of quarks and gluons with all the dynamical effects taken into account. Quark and gluon production from a space-time dependent chromofield is needed to study the production and equilibration of a quark-gluon plasma in ultra relativistic heavy-ion collisions at RHIC and LHC.

We again mention here that the fermion pair production from the space-time dependent field is studied in the literature, but gluon production from space-time dependent chromofield is not studied so far. This is because a consistent theory involving the interaction of the gluons with the classical chromofield is not available in the conventional theory of QCD. The production of $\mathbf{q\bar{q}}$ pairs from a non-abelian field via vacuum polarization is simillar to that of the production of e^-e^+ pairs from the abelian field. This is because the interaction lagrangian of the quantized Dirac field with the classical gauge potential is similar in both the cases.

In high energy collisions, jets are the signatures of quark and gluon production.

For $\mathbf{m_H < 140GeV}$, the most promising discovery mode for the Higgs boson at the LHC has involved the production via gluon fusion, $gg \to H$, followed by the rare decay into two photons, $H \to \gamma\gamma$.

The transverse distribution of particle production from strong constant chromo-electric fields has been explicitly calculated in Ref. 1 for soft-gluon production and in Ref. 2 of ref.[30] for quark (antiquark) production. At high energy large hadron colliders, such as RHIC (Au-Au collisions at $\sqrt{\mathbf{s}} = \mathbf{200GeV}$ [31] and LHC (Pb-Pb collisions at $\sqrt{\mathbf{s}} = \mathbf{5.5TeV}$ [32], about half the total center-of-mass energy, E_{cm}, goes into the production of a semi-classical gluon field, which can be thought to be initially in a Lorentz contracted disc. The gluon field in SU(3) is described by two Casimir invariants, the first one, $\mathbf{C_1 = E^aE^a}$, being related to the energy density of the initial field, whereas the second one, $\mathbf{C_2 = [d_{abc}E^aE^bE^c]}$, is related to the SU(3) color hypercharge left behind by the leading particles.

This was already evident from the Schwinger calculation of the production of a fermion-antifermion pair by an electric field E that is constant in space and time, namely,

$$\frac{\mathbf{dW}}{\mathbf{d^4x}} = \frac{(\mathbf{qeE})^2}{4\pi^2}\sum_{n=1}^{\infty}\mathbf{n^{-2}e^{-n\lambda m^2/(|q|eE)}}$$

where q denotes the charge of the fermion.

This calculation was generalized recently to the nonabelian color group SU(3)c in the special E^a case in which (i) there is only a chromoelectric field, E^a, i.e., the chromomagnetic field $B^a = 0$, (ii) E^a is a constant in space and time, and (iii) all of the group components of E^a point along the same spatial direction.

For example, for $\mathbf{q\bar{q}}$ it was found that

$$\frac{\mathbf{dW_{q\bar{q}}}}{\mathbf{d^4xd^2p_T}} = -\frac{1}{4\pi^3}\sum_{j=1}^{3}|\mathbf{g\lambda_{q,j}}|\mathbf{ln}\left[1-e^{-\pi(p_T^2+m^2)/|g\lambda_{q,j}|}\right] \tag{77}$$

where p_T denotes the momentum of the quark transverse to the direction of the chromoelectric field $\mathrm{E^a} = \mathbf{E^a\hat{z}}$ and where the sums of $SU(3)_c$ group indices a, b, c are from 1 to 8. Integration over p_T yields

$$\frac{\mathbf{dW_{q\bar{q}}}}{\mathbf{d^4x}} = \frac{1}{4\pi^2}\sum_{j=1}^{3}(\mathbf{g\lambda_{q,j}})^2\sum_{n=1}^{\infty}\mathbf{n^{-2}e^{-n\pi m_q^2/|g\lambda_{q,j}|}} \tag{78}$$

where the $\lambda_{q,j}$, depends on three independent gauge and Lorentz invariant eigenvalues λ_j of $f^{abc}E^c$ in SU(3)

$$\lambda_1 = \sqrt{\frac{C_1}{3}}\cos\theta\,,\ \lambda_{2,3} = \sqrt{\frac{C_1}{3}}\cos\left(\frac{2\pi}{3}\pm\theta\right) \tag{79}$$

where

$$\mathbf{C_1(t) = [E^a(t)E^a(t)]}$$

The following result was obtained for the probability of gluon pair production from arbitrary time dependent chromo-electric field $E^a(t)$ in $\alpha = 1$ gauge via Schwinger mechanism[30,33,34]:

$$\mathbf{f(p_T,\theta,C_1)} = \frac{\mathbf{dW_{g(\bar{g})}}}{\mathbf{dtd^3xd^2p_T}} = -\frac{1}{4\pi^3}\sum_{j=1}^{3}|\mathbf{g\lambda_j(t)}|\mathbf{ln}\left[1+e^{-\pi(p_T^2)/|g\lambda_j(t)|}\right] \tag{44}$$

$$\lambda_1^2 = \frac{C_1(t)}{2}[1-\cos\theta(t)]\ \ \lambda_{2,3}^2 = \frac{C_1(t)}{2}\left[1+\cos\left(\frac{\pi}{3}\pm\theta(t)\right)\right]$$

At RHIC and LHC heavy-ion colliders the classical color field play an important role to study production of quark-gluon plasma.

In these situations it is necessary to know how

QCD coupling constant depends on SU(3) color field. In these papers is solved the renormalization group equation in QCD in the presence of SU(3) constant chromo-electric field E^a with arbitrary color index a=1,2,...8. Using background field method in QCD is derived β function from the one loop effective action of quark and gluon in the presence of constant chromo-electric field E^a. Using these two facts is determined the exact dependence of the QCD coupling constant α_s on chromo-electric field E^a in SU(3). Thus, is find the QCD coupling constant which results from figure 1[34]: in the presence of SU(3) chromo-electric field as function θ for fixed values of first Casimir

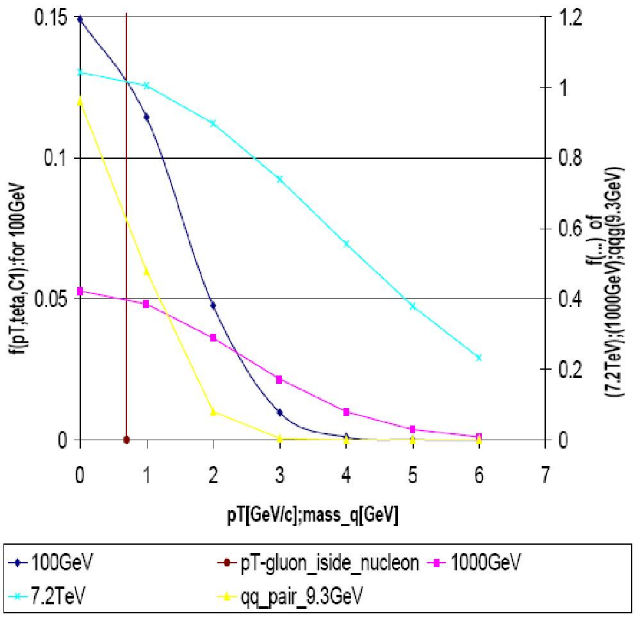

Figure 2 : The evolution of probability of gluon/quarks pairs production for different collision energies

invariant *C*. The λ_j's used are from the above eq. Thus, for $\theta = 0 \to \alpha_s(\lambda_1) \approx 1; \alpha_s(\lambda_2) \approx 0.15; \alpha_s(\lambda_3) = 0.1$, and for $C_1 = 100;1000,7200\text{GeV}^4$, respectively.

$$\alpha_s(\lambda_j) = \frac{g^2(t_j)}{4\pi} = \frac{\alpha_s}{\left[1 + 4\pi\bar{\beta}_0\alpha_s \log\left(\frac{g^2\lambda_j^2}{\mu^4}\right)\right]} = \frac{1}{4\pi\bar{\beta}_0 \log\left(\frac{g^2\lambda_j^2}{\Lambda^4}\right)} \quad (45)$$

$\Lambda = \mu e^{(-1/(4\bar{\beta}_0 g^2))} \cong 200\text{MeV}$; and $0 \le \theta \le 2\pi/3$; $\alpha_s = \frac{g^2}{4\pi}$

$g = 3$

$$\bar{\beta}_0^g = -\frac{11}{32\pi^2}$$

The value of C1 can be estimated from the initial center of-mass energy of the colliding ions, and the volume of the Lorentz contracted Nuclei. For example for gold, $R \approx 10\text{fm}$ and at RHIC the center-of-mass energy is $\approx 200\text{GeV}$ per nucleon. The initial density is then of the order $\rho \cong E_{cm}/(\gamma V_0)$; with $V_0 = 4/3\pi R^3$, and $\gamma = M_{ion}/E_{cm}$. For the above RHIC case $\rho \cong 100\text{GeV}^4$.

The results are plotted in figure 2.

If gluons fragment similar to quarks certainly the fastest gluon jet can be singled out and identified safely like quark jets which start to show up clearly at jet energies of ≈ 3 GeV as had been already shown at SPEAR (Stanford) in 1975[35].

THE CALCULATION OF THE GLUON (MONOPOLE) TRANSVERSE MOMENTUM

In order to point out the value of probability is necessarily to have the value of the gluon (monopole) transverse momentum, for this purpose a model is proposed.

The vector potential of magnetic field produced by magnetic moment of the monopole (gluon) m_{Mo} is

$$A(r) = \frac{\mu_0}{4\pi}\frac{m_{Mo} \times \vec{r}}{r^3}[\text{N/Ampere}] \quad (46)$$

Since, the spin angular momentum, for a monopole $\mu_{Mo} = \frac{gQe\hbar}{2m}$, the color magnetic charge is Q = 68.5, $\vec{s} = \hbar$, $g \cong 2$, and $m_{Mo} = -\frac{\mu_{Mo}S}{\hbar} \cong \mu_{Mo}$, the spin is $S = \hbar/2$, or $m_{Mo} = 2.7e-25[\text{J/Tesla}^{-1}]$

The energy being $T_{P_T} = p_T^2/2m[\text{J}]$, and from the equality of Lorenz force $F_L = q_m(B - vE_g/c^2)[\text{N}]$ with the centripetal force $F_c = mv^2/r$ and with $p_T = mv$, results $p_T = r/v \cdot q_m(B - vE_g/c^2)[\text{N}\cdot\text{s}]$. By using $\tau = \hbar/m_g c^2 = 2.8e-25[\text{s}]$, $m_g \cong 2.2[\text{GeV}]$, result $E_g = A_g/\tau = 2.1e24[\text{N/C}]$, $B = 1.17e15[\text{J/Am}^2]$, and with color magnetic charge $q_m = \frac{2\pi\varepsilon_0\hbar c^2}{q_e}[\text{A}\cdot\text{m}]$, and $r = 0.2[\text{fm}]$, results the transverse momentum for the monopole (gluon) $p_T \cdot c = 0.7[\text{GeV}]$.

With these, as in figure 2, in the case of *pp* collision at LHC, the probability of gluon pair production from arbitrary time $f(p_T, \theta = \pi/3, C_1 = 7.2\text{GTeV}^4) \cong 1$, which in fact that it was happen at LHC.

In the lowest approximation, the Drell-Yan lepton pair of invariant mass M > 1GeV is produced by annihilation of two quarks from the colliding hadrons:

$q_f\bar{q}_f \to \gamma^* \to l^+l^-$

The creation of $e^- - e^+$ pairs in intense laser fields is encountering a growing interest in recent years. It has been stimulated by a pioneering experiment at SLAC (Stanford, USA) where $e^- - e^+$ pair creation was observed in the collision of a 30 GeV γ-photon with an optical laser pulse of $10^{18}\ w/cm^2$. The high-energy photon was first produced by Compton backscattering of the same laser beam off a 46 GeV electron beam.

Due to the high photon density in the intense laser pulse, the simultaneous absorption of more than one laser photon is possible with a non-negligibly small probability. In the experiment, $n = 5$ m laser photons of $\hbar\omega_0 \approx 2[eV]$ combined their energies with the γ-photon upon the collision to overcome the pair creation threshold: $\gamma + n\omega_0 \rightarrow ee^-$ (*nonlinear Breit-Wheeler process*).

The production rate and kinematic distributions of isolated photon pairs produced in hadron interactions are studied[36]. The effects of the initial–state multiple soft–gluon emission to the scattering subprocesses $q\bar{q}$, *qg*, and $gg \rightarrow \gamma\gamma X$ are resummed with the Collins–Soper–Sterman soft gluon resummation formalism.

The effects of fragmentation photons from $qg \rightarrow \gamma q$, followed by $q \rightarrow \gamma X$, are also studied. The results are compared with data from the Fermilab Tevatron collider. A prediction of the production rate and kinematic distributions of the diphoton pair in proton–nucleon reactions is also presented. In this work, the Collins–Soper–Sterman (CSS) soft gluon resummation formalism, developed for Drell–Yan pair (including W and Z boson) production, is extended to describe the production of photon pairs.

In[37] is given a proof of factorization using background field method of QCD.

CONCLUSIONS

We have presented a new analytic approach based on the Dual Ginzburg-Landau theory in order to calculate the strong-field inside nucleons, thus, deriving the values of the monopoles current, the induction, the electromagnetic field, the interaction energies in/between the electric flux tubes as an energy encapsulated by the monopoles circulation, or of the vortex, and between these (giant vortex), respectively.

In the first part we proceed to a review of our analytical model based on the Dual Ginzburg-Landau theory, founding an equivalence with that described in the works from RCNP-Japan, and where, also, is proved the connection between QCD and the dual superconductor scenario.

We provide a detailed analysis of physically important quantities as regarding the nucleons substructure as: the uniform chromoelectric field (vortex) strength inside a nucleon, the mass of monopole viewed as gluons which are the propagators of the QCD and carry colour and anti-colour, with an hedgehog-like configuration, or as a results of interaction of spin-orbit of the monopole current, or of Rashba field interaction, all giving the same result; the quantification of the interaction energies of one vortex ($W^{\pm}$) and of a giant vortex (*GV*), as to be *encapsulated* by the Abrikosov triangular lattice generated by the coalescence of the flux lines.

In the applications, as based on these data, there are calculated: the Higgs boson energy release due of two gluons fusion during the *pp* collision at LHC, gluon pair production from space-time dependent chromofield due of the collision of *pp* and of heavy nuclei;

Due of very promising results in these applications, but mainly of the result of the value of the chromoelectrical field ($8.3\times10^{24}\ V/m$) inside the nucleon, as greater than of Schwinger critical electric field and of parallel magnetic field around the monopoles for $e^+ - e^-$ pairs creation by Schwinger effect, makes possible of *one pair per nucleon* to be obtained. This pair supplies the charges balance (missing) making possible the quarks conversion ($u \rightarrow d$).

Thus, a new understanding of beta decay process is proposed, when a pair of boson $W^- - W^+$ is simultaneously created due of the Schwinger effect when the electrical field (nonabelian) is of maximum value, and near equally with $E_0 \le v.e.v. = E_{cr}^{W^{\pm}}$.

It was discovered that *v.e.v.* is in fact the Schwinger critical field E_{cr} for a pair of $W^{\pm}$ production. This pair decays in beta-electrons which penetrate the condensate barrier by *quantum tunneling* due of the *phase slip* with $2\pi - \varphi$ and of a Φ_0 energy release.

Also, an ad-hoc *bias current* produces a spontaneous suppression of the superconducting order parameter, the model is proved for a free neutron decay. Equally, the same Schwinger pair-production rate is enhanced by a thermal Boltzmann factor in place of quantum tunneling, when this thermalization due of the incidence of an high thermal spike of a photon with valence nucleons destroys the superconductivity.

As a numerical application, is considered the case of ^{26}Al, through its β-decay to 1.809 MeV γ-ray, when at high temperatures ($T_9 = 0.42GK$) equilibrium is reached between ^{26gs}Al and ^{26m}Al which is relevant to some high temperature astrophysical events such as novae, this being proved by our model.

Thus, from the model based on *bias current* it results a necessary photon flux, and from the Schwinger model results the necessary temperature and the duration of the thermal spike; all these parameters can be obtained only by the laser from ELI project (Extreme Light Infrastructure).

APPENDIX A

Massive Vector Boson propagator

In this section, we investigate the propagator of the massive vector boson in the Euclidean metric for the preparation of the analysis on the effective gluon mass in the MA gauge in the Euclidean lattice QCD. We start from the Lagrangian of the free massive vector boson with mass *M* in the Proca formalism following the works of H.Suganuma et al.[7-9]

$$L = \frac{1}{4}(\partial_\mu A_\nu - \partial_\nu A_\mu)^2 + \frac{1}{2}M^2 A_\mu A_\mu \quad (A.1)$$

in the Euclidean metric.

Thus, the scalar-type propagator $G_{\mu\mu}(r;M)$ can be expressed as

$$G_{\mu\mu}(r;M) = \frac{1}{4\pi^2}\frac{M}{r}K_1(Mr) + \frac{1}{M^2}\delta^4(x) \quad (A.2)$$

In the infrared region, the asymptotic expansion

$$K_1(Mr) \cong \sqrt{\frac{\pi}{2Mr}}e^{-Mr}\sum_{n=0}^{\infty}\frac{\Gamma(3/2+n)}{n!\Gamma(3/2-n)}\frac{1}{(2Mr)^n}$$

is applicable for large *Mr*,

$$G^{ch}_{\mu\mu}(r) \equiv \langle A^+_\mu(x)A^-_\mu(y)\rangle = \frac{1}{6}\sum_{a\neq 3,8}\langle A^a_\mu(x)A^a_\mu(y)\rangle$$

$$G^{ch}_{\mu\mu}(r) \equiv \langle A^+_\mu(x)A^-_\mu(y)\rangle \cong A^2_{poloidal} = \frac{1}{6}\left\{\langle A^1_\mu(x)A^1_\mu(y)\rangle + \langle A^2_\mu(x) + A^2_\mu(y)\rangle\right\}$$

Next, we calculate the scalar-type propagator $G_{\mu\mu}(r)$ as the function of the fourdimensional distance $r = \sqrt{(x_\mu - y_\mu)^2}$ in the MA gauge with the U(1)*3*-Landau gauge.

In this gauge, $U_\mu(s)$ is determined without the ambiguity on local gauge transformation.

Here, we study the scalar-type propagator of the diagonal (neutral) gluon as $G^{Abel}_\mu(r) = \langle A^3_\mu(x)A^3_\mu(y)\rangle$ and that of the off-diagonal (charged) gluon as

$$G^{ch}_{\mu\mu}(r) \equiv \langle A^+_\mu A^-_\mu(y)\rangle = \frac{1}{2}\left\{\langle A^1_\mu(x)A^1_\mu(y)\rangle + \langle A^2_\mu(x) + A^2_\mu(y)\rangle\right\}$$

with the charged gluons $A^\pm_\mu = \frac{1}{\sqrt{2}}\{A^1_\mu(x) \pm iA^2_\mu(x)\}$. In $G^{ch}_{\mu\mu}(r)$, the imaginary part, $-\frac{i}{2}\{A^1_\mu(x)A^2_\mu(y) - A^2_\mu(x)A^1_\mu(y)\}$, disappears automatically due to the symmetry

$$\langle A^1_\mu(x)A^2_\mu(y)\rangle = \langle A^2_\mu(x)A^1_\mu(y)\rangle$$

and $(i\partial_\mu \pm eA^3_\mu)A^\pm_\mu = 0$

Therefore, the eq.(A.2) reduces to

$$G^{ch}_{\mu\mu}(r;M) \cong \frac{3M}{2r(2\pi)^{3/2}}\frac{e^{-Mr}}{(Mr)^{1/2}} \quad (A.3)$$

or,

$$rG^{ch}_{\mu\mu}(r) = A^2_{poloidal}x\varepsilon_0 c^2 = (Bx)^2 x\varepsilon_0 c^2[J], \quad A\left[\frac{J}{Am}\right]$$

where $r = x = 2.e-16[m]$, and the "poloidal" nonabelian electric field it could be expressed also as

$$A(r) = \frac{\mu_0}{4\pi}\frac{m \times \vec{r}}{r^3}$$

or of the *Landau gauge*

$$A(x) = Bx\left[\frac{N}{A}\right]; B = 2.6e15\left[\frac{J}{Am^2}\right] from(34),$$

Therefore, (A.3) becomes:

$$2/3\, rG_{\mu\mu}(r;M)(2\pi)^{3/2} \cong M\frac{e^{-Mr}}{(Mr)^{1/2}}$$

At the distance $r \equiv \sqrt{(x_\mu - y_\mu)^2} \cong 2\lambda \cong 0.2fm$, results by trials:

$$6.56e-10[J] \cong M\frac{1}{(Mr)^{1/2}}\exp\left(-\frac{Mr}{\hbar c}\right) =$$

$$M\frac{1}{0.51}\exp\left(-\frac{3.7e-10 \times 2.e-16}{1.e-34 \times 2.997e8}\right) \approx$$

$$1.8M \to M \cong 3.6e-10[J] \Rightarrow 2.25[GeV]$$

and $r_M = \frac{1.e-34 \times 2.997e8}{3.6e-10} \approx 0.1fm \approx \lambda$, just as it was supposed before.

Where the Yukawa-type damping factor e^{-Mr} expresses the short-range interaction in the coordinate space.

In the lattice calculation[4-6] the mass *M* of the vector field $A_\mu(x)$ is estimated from the slope in the logarithmic plot of $\frac{r^{3/2}}{\sqrt{M}}G_{\mu\mu}(r;M)$ as the function of *r*,

$$\ln\left\{\frac{r^{3/2}}{\sqrt{M}}G_{\mu\mu}(r;M)\right\} \cong -Mr + const. \quad (A.4)$$

Next, we calculate the scalar-type propagator $G_{\mu\mu}(r)$ as the function of the fourdimensional distance $r = \sqrt{(x_\mu - y_\mu)^2}$ in the MA gauge with the U(1)*3*-Landau gauge.

In this gauge, $\mathbf{U}_\mu(\mathbf{s})$ is determined without the ambiguity on local gauge transformation.

As shown in Figure 1 of ref.[7], the diagonal-gluon propagator $\mathbf{G}^{\text{Abel}}_{\mu\mu}(\mathbf{r})$ and the charged-gluon propagator $\mathbf{G}^{\text{ch}}_{\mu\mu}(\mathbf{r})$ manifestly differ in the MA gauge. The off-diagonal (charged) gluon $\mathbf{A}^{\pm}_{\mu}$ seems to propagate only within the short range as $\mathbf{r} \le 0.4\mathbf{fm}$, while the diagonal gluon $\mathbf{A}^{3}_{\mu}$ propagates over the long distance. Thus, we find the 'infrared abelian dominance' for the gluon propagator in the MA gauge.

Following the works of Suganuma[4-6] this charged-gluon effective mass $\mathbf{M}_{\text{ch}}$ is considered to be induced by the MA gauge fixing. Due to the effective mass $\mathbf{M}_{\text{ch}} \cong \mathbf{1GeV}$, the charged gluon propagation is restricted within about $\mathbf{M}^{-1}_{\text{ch}} = \mathbf{0.2fm}$. Then, in the infrared region as $\mathbf{r} >> \mathbf{0.2fm}$, the charged gluons $\mathbf{A}^{\pm}_{\mu\mu}$ can not contribute, and only the diagonal gluon $\mathbf{A}^{3}_{\mu}$ can contribute to the long-range physics in the MA gauge. In conclusion, the effective-mass generation of the charged gluon in the MA gauge is considered as the physical origin of the abelian dominance in the infrared region.

Then, the origin of the infrared abelian dominance has been physically explained as the generation of the charged gluon mass $\mathbf{M}_{\text{ch}}$ induced by the MA gauge fixing. On the other hand, in the MA gauge, the charged gluon effects become negligible and the system can be described only by the diagonal gluon component at the long distance as $\mathbf{r} >> \mathbf{M}^{-1}_{\text{ch}} \cong \mathbf{0.2fm}$. For the short distance as $\mathbf{r} \le \mathbf{M}^{-1}_{\text{ch}} \cong \mathbf{0.2fm}$, the effect of charged gluons appears, and hence all the gluon components have to be considered even in the MA gauge.

Alternative ways for monopole mass calculation

The derivation of the Rashba Hamiltonian

The Rashba effect[14] is a direct result of inversion symmetry breaking in the direction perpendicular to the two-dimensional plane. Therefore, let us add to the Hamiltonian a term that breaks this symmetry in the form of an electric field

$$\mathbf{H}_{\text{E}} = -\mathbf{E}_0 \mathbf{z} \tag{A.29}$$

Due to relativistic corrections an electron moving with velocity v in the electric field will experience an effective magnetic field B

$$\mathbf{B} = (\mathbf{v} \times \mathbf{E})/\mathbf{c}^2 \tag{A.30}$$

This magnetic field couples to the electron spin[14]

$$\mathbf{H}_{\text{so}} = \frac{\mu_{\text{B}}}{\mathbf{2c}^2}(\mathbf{v} \times \mathbf{E}) \cdot \sigma \tag{A.31}$$

where the factor 1/2 is a result of the Thomas precession.

The Rashba field E_R exists at the interface and creates the monopole current j_m near the interface. The width of the monopole current distribution, d, is comparable to the decay length of the magnetization at the interface. The monopole current induces the electric current j via Ampe'reʼs law at the interface.

Within this toy model, the Rashba Hamiltonian is given by

$$\mathbf{H}_{\text{R}} = \boldsymbol{\alpha}_{\text{R}}(\boldsymbol{\sigma} \times \mathbf{p}_{\text{T}}) \cdot \hat{\mathbf{z}} \tag{A.32}$$

where $\boldsymbol{\alpha}_{\text{R}} = \frac{\boldsymbol{\mu}_{\text{B}}\mathbf{E}_0}{\mathbf{2mc}^2}$, σ = 1 from Pauli matrix, p_T transverse magnetic moment of the monopole (gluon around); E_0-the electric field induced by an quarks pair $\mathbf{q\bar{q}}$; m-the mass of the monopole. With $\mathbf{p}_{\text{T}} \cdot \mathbf{c} = \mathbf{0.7[GeV]}$ as calculated below, and with $\mathbf{E}_0 = \mathbf{8.33e24[N/C]}$, result: $\boldsymbol{\alpha}_{\text{R}} = \mathbf{3.2e9}$, and $\mathbf{H}_{\text{R}} = \mathbf{1.2e-09[J]}$, respectively.

The magnetic moment of the electron is

$$\mathbf{m}_{\text{s}} = -\frac{\mathbf{g}_{\text{s}}\boldsymbol{\mu}_{\text{B}}\mathbf{S}}{\hbar} \tag{A.33}$$

where $\mu_{\text{B}} = \mathbf{9.27 \times 10^{-24}[JT^{-1}]}$, μ_{B} is the Bohr magneton, $\mathbf{S} = \hbar/2$ is electron spin, and the g-factor g_s is 2 according to Dirac's theory, but due to quantum electrodynamic effects it is slightly larger in reality: 2.002, for a muon g = 2.

The Bohr magneton is defined in *Si* units by

$$\boldsymbol{\mu}_{\text{B}} = \frac{\mathbf{e}\hbar}{\mathbf{2m}_{\text{e}}} \tag{A.34}$$

The vector potential of magnetic field produced by magnetic moment m_{Mo} is

$$\mathbf{A(r)} = \frac{\boldsymbol{\mu}_0}{\mathbf{4\pi}} \frac{\mathbf{m}_{\text{Mo}} \times \vec{\mathbf{r}}}{\mathbf{r}^3} \tag{A.35}$$

and magnetic flux density is

$$\mathbf{B(r)} = \nabla \times \mathbf{A} = \frac{\boldsymbol{\mu}_0}{\mathbf{4\pi}}\left(\frac{\mathbf{3}\vec{\mathbf{r}}(\mathbf{m}_{\text{Mo}} \cdot \vec{\mathbf{r}})}{\mathbf{r}^5} - \frac{\mathbf{m}_{\text{Mo}}}{\mathbf{r}^3}\right) \cong \mathbf{A/r} \tag{A.36}$$

$$\mathbf{j}_{\text{e}} = -\frac{\mathbf{c}}{\mathbf{4\pi}}\frac{\mathbf{dB}}{\mathbf{dr}}(\mathbf{4\pi c}\boldsymbol{\varepsilon}_0) = -\frac{\mathbf{c}^2\boldsymbol{\varepsilon}_0\mathbf{A}}{\mathbf{r}^2} \tag{A.37}$$

We can calculate the observable spin magnetic moment (a vector), $\vec{\boldsymbol{\mu}}_{\text{S}}$, **for a sub-atomic particle with charge *q*, mass *m*, and spin angular momentum (also a vector), $\vec{\mathbf{s}}$, via:**

$$\vec{\boldsymbol{\mu}}_{\mathbf{s}} = \mathbf{g}\frac{\mathbf{q}}{\mathbf{2m}}\vec{\mathbf{s}}$$

Therefore, for a monopole

$$\mu_{Mo} = \frac{gQe\hbar}{2m_{Mo}}, \; Q = 68.5, \; \vec{s} = \hbar, \; g \cong 2 \quad (A.38)$$

And

$$m_{Mo} = -\frac{\mu_{Mo} S}{\hbar} \cong \mu_{Mo}$$

(A.39)

Numerically, results: $\mu_{Mo} = 2.7e-25[J/T]$, $A = 0.6[N/A]$,

$$B = A/r = 2.86e15[J/Am^2]$$

(A.40)

if we consider $r = 0.2\ fm$ results

$H_{So} = 3.9e-10[J] \rightarrow M_{monopole} = 2.45[GeV]$, and the electric current is $j_e = 1.2e7[A/fm^2]$ which is comparable with the magnetic current

$j_\varphi = 1.15e7[A/fm^2]$, as from eq. (35.2).

The Rashba interaction contributes to the DC monopole current at the interface[14].

REFERENCES

[1] G.'t Hooft; in High Energy Physics, A.Zichici, (Ed); Editrice Compositori, Bologna, (1976).

[2] (a) S.Mandelstam; Phys.Rept.C, 23, 245 (1976); (b) V.Singh, D.A.Browne, R.W.Haymaker; Structure of Abrikosov Vortices in SU(2) Lattice Gauge Theory, Phys.Lett.B, 306, 115 (1993).

[3] Stefan Mehedinteanu; On the Photonuclear rates calculation for nuclear transmutation of beta decay nuclides by application of the Ginzburg-Landau theory, Journal of Nuclear and Particle Physics, 2(3), 57-70 (2012).

[4] H.Suganuma, S.Sasakib, H.Tokib, H.Ichie; Dual Higgs Mechanism for Quarks in Hadrons, (1995).

[5] H.Ichie, H.Suganuma, H.Toki; The System of Multi Color-flux-tubes in the Dual Ginzburg-Landau Theory, (1996).

[6] Y.Koma, H.Suganuma, H.Toki; Flux-Tube Ring and Glueball Properties in the Dual Ginzburg-Landau Theory, (1999).

[7] H.Ichie, H.Suganuma; Semi-analytical Proof of Abelian Dominance on Confinement in the Maximally Abelian Gauge, (1998).

[8] Shinya Gongyoa, Takumi Iritania, Hideo Suganuma; Off-diagonal Gluon Mass Generation and Infrared Abelian Dominance in Maximally Abelian Gauge in SU(3) Lattice QCD, Phys.Rev.D, 86, 094018 (2012).

[9] M.A.Gauge, H.Suganumaa, K.Amemiyaa, H.Ichiea; Infrared Abelian Dominance and Dual Higgs Mechanism, (2004).

[10] H.Suganuma, K.Amemiya, H.Ichie, Y.Koma; Construction of the Dual Ginzburg-Landau Theory from the Lattice QCD, (2008).

[11] M.N.Chernodub, M.I.Polikarpov; Abelian Projections and Monopoles, (1997).

[12] (a) K.H.Bennemann, J.B.Ketterson; Superconductivity conventional and unconventional Superconductors, (chapter J.Pitaevscki), Springer, 1, (2008); (b) Gunnar S.Bali, Christoph Schlichter, Klaus Schilling; Probing the QCD Vacuum with Static Sources in Maximal Abelian Projection, (1998).

[13] N.S.Craigie, et al.; Theory and detection of magnetic monopoles in gauge theories: A collected..., Word Scentific Publishing Co., 2.

[14] Akihito Takeuchi, Gen Tatara; Spin damping monopole, Journal of the Physical Society of Japan, 81, 033705 (2012).

[15] Andrey Chaves, F.M.Peeters, G.A.Farias, M.V.Miloševiæ; Vortex-vortex interaction in bulk superconductors: Ginzburg-Landau theory, Phys.Rev.B, 83, 054516 (2011).

[16] R.Ruffini, Gregory Vereshchagin; Electron-positron pairs in physics and astrophysics: From heavy nuclei to black holes, (2009).

[17] Sang Pyo Kim, Don N.Page; Schwinger Pair Production in Electric and Magnetic Fields, (2006).

[18] S.Mehedinteanu; On the numerical analysis of decay rate enhancement in metallic environment, Acta Physica Polonica B No.10, 38, (2007).

[19] A.Benedetti, R.Ruffini, G.V.Vereshchagin; Phase space evolution of pairs created in strong electric fields, (2013).

[20] Sang Pyo Kim, Hyun Kyu Le; Schwinger Pair Production at Finite Temperature in Scalar QED, (2007).

[21] Carsten Müller; Nonlinear Bethe-Heitler pair creation with attosecond laser pulses at the LHC, Physics Letters B, 672(1), 56-60 (2009).

[22] Tim-Oliver Müller, Carsten Müller; Spin correlations in nonperturbative electron-positron pair creation by petawatt laser pulses colliding with a TeV proton beam, (2010).

[23] Anis Dadi, Carsten Müller; Laser-assisted nuclear photoeffect, (2012).

[24] K.-M.Spohr, S.D.Pain, K.W.D.Ledingham, S.Brien, R.Chapman, M.Hassan, J.Melone, J.F.Smith; Studies of enhanced decay of ^{26}Al in hot plasma environments, www.nanolink.net/e-newsletter...e.../ELI-NP-WhiteBook.pdf.

[25] Alain Coc, Marie-Geneviéve Porquet, Frédéric Nowacki; Lifetimes of ^{26}Al and ^{34}Cl in an Astrophysical Plasma, (1999).

[26] R.Ward, W.Fowler; Thermalization of long-lived nuclear isomers states under Stellar conditions, Astrophysical Journal, 238, 266-286 (1980).

[27] http://www.extreme-light-infrastructure.eu

[28] The White Book of ELI Nuclear Physics Bucharest-Magurele, Romania, The ELI-Nuclear Physics working groups, www.eli-np.ro/documents/ELI-NP-WhiteBook.pdf.

[29] Gouranga C.Nayak, Walter Greiner; Pair Production from Space-Time Dependent Chromofield, (2000).

[30] Gouranga C.Nayak, Robert Shrock; Gauge-Invariant Quantities Characterizing Gauge Fields in Chromodynamics, (2007).

[31] Victor D.Efros, Winfried Leidemann, Giuseppina Orlandini, Edward L.Tomusiak; Photodisintegration of three-body nuclei with realistic 2N and 3N forces, Physics Letters B, http://arxiv.org/abs/nucl-th/9911049v2; 484(3-4), July (2000).

[32] P.G.Thirolfa, L.Csigea, M.Güntherb, D.Habsa; A.Krasznahorkay; Perspectives for spectroscopy of actinides with highly brilliant γ beams, Acta Physica Polonica B, 43(2), (2012).

[33] Y.M.Cho, J.H.Kim; QCD effective action with a most general homogeneous field background, (2007).

[34] G.C.Nayak; Renormalization Group Equation and QCD Coupling Constant in the Presence of SU(3) Chromo-Electric Field, (2009).

[35] Bruno R.Stella, Hans-Jürgen Meyer; (9.46 GeV) and the gluon discovery (a critical recollection of PLUTO results).

[36] C.Balázs, E.L.Berger, S.Mrenna, C.P.Yuan; Photon Pair Production with Soft Gluon Resummation In Hadronic Interactions, (1997).

[37] Gouranga C.Nayak; Proof of Factorization Using Background Field Method of QCD, (2009).