COLOR CONFINEMENT AND DUAL SUPERCONDUCTIVITY: AN UPDATE

ADRIANO DI GIACOMO
Dipartimento di Fisica, Università di Pisa and
INFN Sezione di Pisa
Via Buonarroti 2, I-56127 Pisa, Italy

Abstract. The evidence for dual superconductivity as a mechanism for color confinement is reviewed. New developments are presented for full QCD, i.e. in the presence of dynamical quarks.

1. Introduction

Confinement is by definition the absence of free colored particles in nature. In spite of the clear evidence coming from high energy experiments that quarks and gluons are the fundamental constituents of hadrons, none of them has ever been detected.

Only upper limits exist to production cross sections. The cross section \( \sigma_q \) for the inclusive process

\[
p + p \rightarrow q(\bar{q}) + X
\]

has an upper limit\[1\]

\[
\sigma_q < 10^{-40} \text{cm}^2
\]

(2)

At the same energies the total cross section \( \sigma_T \) has the value

\[
\sigma_T \simeq 10^{-25} \text{cm}^2
\]

(3)

In perturbation theory the ratio \( \sigma_q/\sigma_T \) is expected to be a sizable fraction of unity. From (2) and (3)

\[
\sigma_q/\sigma_T < 10^{-15}
\]

(4)

Relic quarks in nature have been hunted since they were first proposed as fundamental bricks of matter\[2\].
Fourty years of Millikan-like experiments looking for fractionally charged particles have produced as upper limit \[1\]

\[
n_q/n_p < 10^{-27}
\]  

(5)

for the ratio of the abundance of quarks \(n_q\) and that of nucleons \(n_p\) in nature. The limit (5) results from the analysis of \(\sim 1\)g of matter and no quarks found. In the absence of confinement the Standard Cosmological Model predicts\[3\]

\[
n_q/n_p \simeq 10^{-12}
\]  

(6)

Again

\[
\frac{(n_q/n_p)_{\text{obs}}}{(n_q/n_p)_{\text{expected}}} < 10^{-15}
\]  

(7)

The only natural explanation of the limits (3) and (7) is that \(\sigma_q, n_q\) are exactly zero, or that confinement is an absolute property based on symmetry \[4\]. Similar situations are e.g.

i) the photon mass, which is experimentally bounded by the inverse radius of the solar system\[1\]. The corresponding symmetry is gauge invariance.

ii) ordinary superconductivity. The upper limit of the resistivity of superconductors is many orders of magnitude smaller than that of any other material. The symmetry pattern behind that is the Higgs phenomenon\[5\].

2. Virtual tests (lattice)

QCD at finite temperature can be simulated on the lattice. It is a well known theorem that the partition function of a field system is equal to the euclidean Feynman integral, with immaginatory time ranging from 0 to \(1/T\), and periodic boundary conditions for bosons, antiperiodic for fermions:

\[
Z = \int [\mathcal{D}\varphi] \exp \left[ -\int_0^{1/T} dt \int d^3x \mathcal{L}(\vec{x}, t) \right]
\]  

(8)

If the lattice size in the time direction is \(N_T\) then

\[
T = \frac{1}{N_T a}
\]  

(9)

\(a\) being the lattice spacing in physical units. Renormalization group gives, at large enough \(\beta = 2N/g^2\), \(a = \frac{1}{N_L} \exp(-\beta/b_0)\), where \(b_0 > 0\) by asymptotic freedom, i.e.

\[
T = \frac{\Lambda_L}{N_T} \exp(\beta/b_0)
\]  

(10)
Low T corresponds to large $g^2$ (strong coupling or disorder in the language of statistical mechanics); high T to weak coupling or order. This is the opposite to what happens in ordinary spin systems, for which $T$ plays the role of $g^2$.

In pure gauge theories an order parameter for confinement is $\langle L \rangle$, the Polyakov line, which is the trace of the parallel transport along the imaginary time axis from 0 to $1/T$, closed by periodic boundary conditions. The corresponding symmetry is $Z_N$.

On the lattice the correlator

$$G(\vec{r}) = \langle L(\vec{r})L^\dagger(\vec{0}) \rangle$$

is measured [6]. Cluster property requires

$$G(\vec{r})_{r \to \infty} \sim A \exp(-\sigma a N_T r) + |\langle L \rangle|^2$$

The potential energy of a static $q\bar{q}$ pair at distance $r$ is given by

$$V(r) = -\frac{1}{a N_T} \ln G(r)$$

A temperature $T_c$ is found such that for $T < T_c$ $\langle L \rangle = 0$ or

$$V(r)_{r \to \infty} \sim \sigma r$$

which means confinement. For $T > T_c$ $\langle L \rangle \neq 0$ and

$$V(r)_{r \to \infty} \sim \text{const.}$$

which means deconfinement.

Finite size scaling analysis of the correlator around the critical point provides a determination of the critical index $\nu$. For $SU(2)$ pure gauge theory the transition is second order [6], consistent with the class of universality of the 3d ising model ($\nu = .62$) as expected [7], and $T_c/\sqrt{\sigma} \simeq .7$. For $SU(3)$ pure gauge theory the transition is weak first order [8, 9], ($\nu = .33$) and $T_c/\sqrt{\sigma} \simeq .65$, which, by the usual assumption $\sqrt{\sigma} = 425$ MeV gives $T_c \simeq 270$ MeV.

In the presence of dynamical quarks $Z_N$ is explicitly broken, and $\langle L \rangle$ cannot be an order parameter. For two equal-mass dynamical quarks the situation is depicted in fig.1.

The transition temperature is determined, at given quark mass, by looking at the maximum of a number of susceptibilities, e.g. $\int (\bar{\psi}\psi(x)\bar{\psi}\psi(0)) d^3x$, $\int \langle L(x)L(0) \rangle d^3x$. All of them show a maximum at the same $T_c$ [10]. For high enough $m_q$, ($m_q \geq 3$ GeV) the maximum of the Polyakov line susceptibility goes large with increasing spatial volume as in the quenched case: a finite
size scaling analysis shows that the transition is first order. There are indications that the transition is second order in the chiral limit \( m_q = 0 \), as suggested by symmetry arguments\[11\]. At intermediate values of \( m_q \) none of the susceptibilities which have been considered increases with increasing spatial volume, and a possible conclusion is that there is no phase transition but only a crossover. The overall situation is rather confusing. It is not clear a priori what is the relation between chiral symmetry and confinement. It is not fully clear either what susceptibilities are entitled to determine the order of the transition by their behavior at large volumes. In principle the relevant quantities should be those appearing in the expression of the free energy. The free energy (effective lagrangean ) depends on the dominant excitations and on their symmetry. What are the dominant excitations is exactly the problem under investigation.

3. Duality.

Confined phase is disordered. How can the symmetry of a disordered phase be defined? The key concept is duality\[12\]. It applies to \( d \)-dimensional systems admitting non trivial topological excitations in \((d-1)\) dimensions. These systems admit two complementary descriptions.

1) A direct description in terms of the fields \( \phi \), with order parameters \( \langle \phi \rangle \), in which the topological configurations \( \mu \) are non local. This description
is convenient in the weak coupling regime \((g \ll 1)\), i.e. in the ordered phase.

2) A dual description in which the topological excitations \(\mu\) become local fields, and the original fields \(\phi\) topological configurations. The dual coupling \(g_D\) is related to \(g\) as \(g_D \sim 1/g\). This description is convenient in the disordered phase (strong coupling regime). Its symmetry is described by \(\langle \mu \rangle\) (disorder parameter). Duality maps the strong coupling regime of the direct description into the weak coupling regime of the dual description.

The prototype system for duality is the Ising model\(^{[13]}\) where dual excitations are kinks. Other examples are \(\text{N}=2\) SUSY QCD\(^{[14]}\), where the dual excitations are monopoles; M string theories\(^{[15]}\); 3-d XY model, where dual excitations are abelian vortices\(^{[16]}\); 3-d Heisenberg magnet, with 2-d Weiss domains as dual excitations\(^{[17]}\); compact U(1) gauge theory, where dual excitations are monopoles\(^{[18, 19]}\).

In QCD the dual topological excitations have to be identified: as we will see, however, information exists on their symmetry. Two original proposals exist in the literature, which have been widely studied:

a) Monopoles\(^{[20, 21]}\). The idea is that vacuum acts as a dual superconductor, which confines electric charges by Meissner effect, in the same way as magnetic charges are confined in an ordinary superconductor.

Developments of this approach will be the subject of the next sections.

b) Vortices\(^{[4]}\). The symmetry involved is \(Z_N\).

In 2+1 dimensions a conserved charge exists, the number of vortices minus the number of antivortices, and vortices are described by a local field. In 3+1 dimensions a dual Wilson loop can be defined ('tHooft loop) \(B(C)\), in connection with any closed path \(C\). The algebra which is obeyed by \(B(C)\) and by the ordinary Wilson Loop \(W(C')\) is

\[
B(C)W(C') = W(C')B(C) \exp \left( \frac{2\pi i n_{CC'}}{N_c} \right) \tag{16}
\]

where \(n_{CC'}\) is the linking number of the two loops. From eq(16) it follows that, if \(\langle W(C') \rangle\) obeys the area law \(\langle B(C) \rangle\) obeys the perimeter law, and if \(\langle B(C) \rangle\) obeys the area law then \(\langle W(C') \rangle\) obeys the perimeter law. If we denote by \(\langle L \rangle\) the ordinary Wilson loop which wraps the lattice through periodic b.c. in time (Polyakov loop), and by \(\langle \tilde{L} \rangle\) the analogous dual loop (‘tHooft’s line), then in the confined phase \(\langle L \rangle = 0, \langle \tilde{L} \rangle \neq 0\), whilst in the deconfined phase \(\langle \tilde{L} \rangle = 0, \langle L \rangle \neq 0\). \(\langle \tilde{L} \rangle\) is a disorder parameter for confinement. These relations have been tested on the lattice\(^{[22, 23]}\). The corresponding symmetries \(Z_N\) and \(\tilde{Z}_N\) are explicitly broken in the presence of fermions.
4. Monopoles.

Monopoles in non abelian gauge theories are always abelian (Dirac) monopoles. This statement can be immediately checked by looking at the field produced by a static configuration of colored matter at large distances, by use of the familiar multipole expansion [24]. Monopoles are identified by a constant diagonal matrix in the algebra, with integer or half-integer values: they carry N-1 abelian magnetic charges. The same physics emerges from the procedure known as abelian projection [21]. We shall illustrate it for SU(2): the general case [25] is not substantially different. Let \( \varphi(x) \) be any operator in the adjoint representation, and \( \bar{\varphi}(x) = \varphi(x)/|\varphi(x)| \) its direction in color space. Define [26]

\[
F_{\mu\nu} = \bar{\varphi} G_{\mu\nu} - \frac{1}{g} \bar{\varphi} (D_\mu \bar{\varphi} \wedge D_\nu \bar{\varphi})
\]

with \( G_{\mu\nu} = \partial_\mu \tilde{A}_\nu - \partial_\nu \tilde{A}_\mu + g \tilde{A}_\mu \wedge \tilde{A}_\nu \) the field strength and \( D_\mu = \partial_\mu + g \tilde{A}_\mu \wedge \) the covariant derivative. Both terms in eq.(17) are color singlets and gauge invariant; the combination is chosen to cancel bilinear terms \( A_\mu A_\nu \). Indeed one has identically:

\[
F_{\mu\nu} = \bar{\varphi} (\partial_\mu \tilde{A}_\nu - \partial_\nu \tilde{A}_\mu) - \frac{1}{g} \bar{\varphi} (\partial_\mu \bar{\varphi} \wedge \partial_\nu \bar{\varphi})
\]

In a gauge in which \( \bar{\varphi} \) is constant, e.g. \( \bar{\varphi} = (0,0,1) \), \( F_{\mu\nu} \) is abelian:

\[
F_{\mu\nu} = \partial_\mu A_3^\nu - \partial_\nu A_3^\mu
\]

A magnetic current, \( j_\mu \), can be defined in terms of the dual tensor \( F^*_{\mu\nu} = \frac{1}{2} \varepsilon_{\mu\nu\rho\sigma} F_{\rho\sigma} \).

\[
j_\mu = \partial_\nu F^*_{\mu\nu}
\]

\( j_\mu \) is identically zero (Bianchi identities) in a non compact formulation of the theory. In a compact formulation, like Lattice, \( j^\mu \) can be non zero. In any case it is identically conserved

\[
\partial_\mu j^\mu = 0
\]

Magnetic charges are Dirac monopoles, obeying Dirac quantization condition \( Q = n/2g \). The corresponding magnetic \( U(1) \) symmetry can either be realized a la Wigner, and then the Hilbert space consists of superselected sectors with definite magnetic charge, or Higgs-broken, and then the system behaves as a dual superconductor. If the ideas of ref’s [20, 21] are correct the expectation is that QCD vacuum behaves as a dual superconductor (Higgs-broken phase) for \( T < T_c \), and as a magnetic superselected system for \( T > T_c \). A disorder parameter should discriminate between superconductor
and normal. Such a parameter has been constructed \cite{27, 28, 29} as the v.e.v. \langle \mu \rangle of an operator \mu carrying magnetic charge. In fact \mu is a Dirac-like operator \cite{30}, charged and gauge invariant \cite{18, 33}. The construction of \mu is at the level of a theorem for compact U(1) \cite{18, 33}. In non abelian gauge theories it is undefined by terms \mathcal{O}(a^2), a being the lattice spacing, like the abelian projection itself \cite{32, 34}.

5. Results

The basic structure of \mu is a translation of the field configuration in the Schrödinger picture by a classical monopole configuration. In the same way as

\[ e^{ipa}|q\rangle = |q + a\rangle \]  

defining

\[ \mu(\vec{x}, t) = \exp \left( i \int d^3 \vec{y} \Pi(\vec{y}, t) \vec{\varphi}(\vec{x} - \vec{y}) \right) \]

with \Pi(\vec{y}, t) the conjugate momentum to the field \varphi(\vec{y}, t), and \vec{\varphi}(\vec{y}, t) the classical field configuration to be added

\[ \mu(\vec{x}, t)|\varphi(\vec{y}, t)\rangle = |\varphi(\vec{y}, t) + \vec{\varphi}(\vec{x} - \vec{y})\rangle \]  

In fact the basic structure has to be adapted to a compact formulation, in which the field cannot be translated at will \cite{19}, and to the abelian projected situation, in which only the abelian part of the field has to be translated. All this has been done \cite{27}. The resulting disorder parameter \langle \mu \rangle can be finally written as the ratio of two partition functions

\[ \langle \mu \rangle = \frac{\tilde{Z}(\beta)}{Z(\beta)} \]  

with \[ Z(\beta) = \int [\mathcal{D}\varphi] \exp(-\beta S), \] and

\[ Z(\beta) = \int [\mathcal{D}\varphi] \exp \left[ -\beta (S + \Delta S) \right] \]  

\( S + \Delta S \) is obtained from \( S \) by a modification of the space time plaquettes at time \( t, \Pi_{0i}(\vec{n}, t) \); for details see ref's \cite{19, 27}. Instead of \( \langle \mu \rangle \) itself it is more convenient to study

\[ \rho = \frac{d}{d\beta} \ln(\mu) = \langle S \rangle_S - \langle S + \Delta S \rangle_S \Delta S \]  

On the one hand \( \rho \) is numerically easier, since \( \langle \mu \rangle \) fluctuates wildly as any partition function; on the other hand we shall see that \( \rho \) contains all the
relevant information as $\langle \mu \rangle$, and also in a more convenient way. From $\rho$, $\langle \mu \rangle$ is obtained as

$$\langle \mu \rangle = \exp \left[ \int_{0}^{\beta} \rho(x) dx \right]$$

(25)

since $Z(\beta = 0) = \tilde{Z}(\beta = 0) = 1$.

The typical shapes of $\langle \mu \rangle$ and $\rho$ as functions of $\beta$ are plotted in fig.2. The position of the negative peak coincides with the deconfining phase transition.

![Figure 2. Typical shape of $\langle \mu \rangle$ and $\rho$.](image)

Fig.3 shows $\rho$ for different spatial sizes $N_s$ of the lattice, at fixed $N_t = 4$, for $SU(2)$ pure gauge theory.

The position of the peak coincides with the maximum of the susceptibility of the Polyakov line, as determined in ref[4], i.e. with the phase transition. In the range of temperatures $T < T_c$ $\rho$ stays practically constant by increasing the volume, as shown in fig.4, and this means that $\langle \mu \rangle$ has a non zero limit in that region.

In the range $T > T_c$ $\rho$ diverges to $-\infty$ as $N_s$ goes large, as

$$\rho = -kN_s + k' \quad k > 0$$

(26)
as shown in fig.5.

If we had measured $\langle \mu \rangle$ itself instead of $\rho$ we would have found it consistent with zero within large errors at large $N_s$. Measuring instead $\rho$ and checking numerically the behavior eq.(31) amounts to state that $\langle \mu \rangle$ is exactly zero in the thermodynamical limit. In the critical region, around $T_c$ one expects

$$\langle \mu \rangle \overset{T \to T_c}{\sim} \tau^\delta \Phi \left( \frac{a}{\xi}, \frac{N_s a}{\xi} \right)$$

(27)

with $\tau = 1 - \frac{T}{T_c} \propto (\beta_c - \beta)$, $\delta$ a critical index, $a$ the lattice spacing and $\xi$ the correlation length. The transition is known to be second order for $SU(2)$,
weak first order for $SU(3)$ gauge theory, and therefore $\xi \sim \tau^{-\nu}$ goes large at small $\tau$’s. Neglecting $a/\xi \sim 0$ and trading the variable $N_s/\xi$ with $\tau N_s^{1/\nu}$, the scaling law follows for $\rho$ from eq.(32)

$$\frac{\rho}{N_s^{1/\nu}} = \frac{\delta}{\tau N_s^{1/\nu}} + \Phi'(0, \tau N_s^{1/\nu})$$

(28)

For different sizes of the lattice $\rho/N_s^{1/\nu}$, when plotted versus $\tau N_s^{1/\nu}$ is expected to be independent of $N_s$. A best fit to the data allows a determination of $\beta_c, \nu, \delta$. The quality of the scaling is shown in fig.6 for $SU(2)$. The result is

$$SU(2) \quad \nu = .62(1) \quad \delta = .20(3)$$
$$SU(3) \quad \nu = .33(1) \quad \delta = .50(3)$$

All that is for quenched theory. A few different abelian projections have been tested, and the behavior of $\rho$, as well as the value of $\nu$ and $\delta$ are independent of the abelian projection. An additional test has been made with no abelian projection, i.e. by assuming as diagonal operators the nominal $\lambda_3, \lambda_8$ used in the simulation. This amounts to perform an average over all abelian projections, and the result does not change.

We can thus state that the confining phase of a pure gauge theory is a dual superconductor in all the abelian projections, and undergoes a
transition to normal at $T_c$. What we learn from that is that, whatever the dual excitations are, they carry magnetic charge in all the abelian projections. The definition of the operator $\mu$ can be easily extended to full QCD, with dynamical quarks. A natural question is whether also in this case the confining phase is characterized by dual superconductivity. This
Figure 6. Finite size scaling of $\rho$

would indeed be the expectation in the spirit of the $N_c \to \infty$ limit. If the physics of gauge theories is determined by the limit $N_c \to \infty$ at $g^2 N_c$ fixed, corrections $1/N_c$ being small, then the mechanism of confinement has to be $N_c$ independent, and also insensitive to the presence of dynamical quarks, their contribution being non leading in the $1/N_c$ expansion. The parameter $\rho = \frac{d}{d\beta} \ln \langle \mu \rangle$ should go to $-\infty$ for $T > T_c$ in the thermodynamical limit so that $\langle \mu \rangle = 0$: this is indeed the case \cite{35}. For $T < T_c$, $\rho$ converges to a finite limit, i.e. $\langle \mu \rangle \neq 0$ and there is superconductivity \cite{35}. The shaded area of fig.1 does indeed correspond to a superconductor, the upper area to a superselected magnetic system: the negative peak is at the transition as defined in sect.2 [fig.7]

Around $T_c$ a finite size scaling analysis can give information on the order of the transition: the difference with respect to the quenched case is that now an extra dimensionful quantity, the quark mass, enters and there is a two variable scaling. As usual we expect

$$\langle \mu \rangle \sim \tau^\delta \Phi(\frac{a}{\xi}, \frac{N_s a}{\xi}, m_q N_s^\alpha)$$

(29)

In the critical regime $a/\xi \simeq 0$, $N_T/N_s \simeq 0$, $\xi \sim \tau^{-\nu}$ and we can again trade $\xi/N_s$ with $\tau N_s^{1/\nu}$, and the scaling law becomes

$$\langle \mu \rangle \sim \tau^\delta \Phi(0, \frac{N_s a}{\xi}, m_q N_s^\alpha)$$

(30)

The index $\alpha$ is known: one can chose for different values of $N_s$ suitable quark masses, so as to keep $m_q N_s^\alpha$ fixed. The scaling law is then the same
as eq (28), whence the index $\nu$ can be extracted and with it information on the order of the transition. This is a heavy numerical program which is on the way on our APE machines. The result will possibly be relevant to the problems discussed in sect 2.

6. Conclusions and outlook.

1) Confinement is characterized by dual superconductivity of the vacuum in all the abelian projections, both in quenched and in full QCD, in line with the idea of the limit $N_c \to \infty$. A universal disorder parameter has been defined for it.

2) We do not know the dual excitations of QCD, nor their effective lagrangean. However we know that they carry magnetic charge in all the abelian projections.

3) An analysis of the critical region in full QCD is on the way, which will possibly clarify the nature of the phase transition depicted in fig.1.

4) The determination of the parameters of the dual supercondutor, Higgs mass, penetration depth... is fundamental and relevant to understand the deconfinement signals which could come from heavy ion collisions: this is also part of our research program.

The contributions of L. Del Debbio, M. D’Elia, B. Lucini, G. Paffuti to the research program are acknowledged.
References

1. Review of Particle Physics E.P.J. 15 (2000)
2. M. Gellmann, Phys. Rev. 125 (1962) 1067 ;
3. L. B. Okun, Leptons and Quarks, North Holland (1982);
4. G. ’t Hooft, Nucl. Phys. B 138 (1978) 1
5. S. Weinberg, Progr. Theor. Phys. Suppl. 86 (1986) 43;
6. J. Engels, F. Karsch, H. Satz, I. Montvay, Nucl. Phys. B 205 (1982) 239 ;
7. B. Svetitsky, L.G. Yaffe, Nucl. Phys. B 210 (1982) 423;
8. F.R. Brown, N.H. Christ, Y.F. Deng, M.S. Gao, T.J. Woch, Phys. Rev. Lett. 61 (1988) 2058;
9. M. Fukujita, M. Okawa, A. Ukawa, Phys. Rev. Lett. 63 (1989) 1768 ;
10. F. Karsch, E. Laermann, Phys. Rev. D 50 (1994) 6954;
11. R.D. Pisarski, F. Wilczek, Phys. Rev. D 29 (1984) 338
12. H.A. Kramers, G.H. Wannier, Phys. Rev. 60 (1941) 252;
13. L.P. Kadanoff, H. Ceva, Phys. Rev. B 3 (1971) 3918;
14. N. Seiberg, E. Witten, Nucl. Phys. B 341 (1994) 484;
15. T. Banks, W. Fischler, S.H. Shenker, L. Susskind, Phys. Rev. D 55 (1997) 5112
16. G. Di Cecio, A. Di Giacomo, G. Paffuti, M. Triggiante, Nucl. Phys. B 489 (1997) 739;
17. A. Di Giacomo, D. Martelli, G. Paffuti, Phys. Rev. D 60 (1999) 094511;
18. J. Froelich, P.A. Marchetti, Commun. Math. Phys. 112 (1987) 343;
19. A. Di Giacomo, G. Paffuti, Phys. Rev. D 56 (1997) 6816;
20. G. ’t Hooft, High Energy Physics EPS International Conference Palermo 1975, A.Zichichi ed., S. Mandelstam Phys. Rep. 23C (1976) 245;
21. G. ’t Hooft, Nucl. Phys. B 190 (1981) 455;
22. L. Del Debbio, A. Di Giacomo, B. Lucini, Nucl. Phys. B 594 (2001) 287;
23. L. Del Debbio, A. Di Giacomo, B. Lucini, Phys. Lett. B 500 (2001) 326;
24. S. Coleman, Erice Summer School 1981 Plenum Press, A. Zichichi ed.;
25. L. Del Debbio, A. Di Giacomo, B. Lucini, G. Paffuti, hep-lat 0203024;
26. G. ’t Hooft, Nucl. Phys. B 79 (1974) 276
27. A. Di Giacomo, B. Lucini, L. Montesi, G. Paffuti, Phys. Rev. D 61 (2000) 034503;
28. A. Di Giacomo, B. Lucini, L. Montesi, G. Paffuti, Phys. Rev. D 61 (2000) 034504 ;
29. J.M. Carmona, M. D’Elia, A. Di Giacomo, B. Lucini, G. Paffuti, Phys. Rev. D 64 (2001) 114507;
30. P.A.M. Dirac Canad. J. Phys. 33 (1955) 650;
31. J.M. Carmona, M. D’Elia, L. Del Debbio, A. Di Giacomo, B. Lucini, G. Paffuti, Nucl. Phys. Proc. Suppl. B 106 (2002) 607 and in preparation;
32. A. Di Giacomo, G. Paffuti, Nucl. Phys. Proc. Suppl. 106 (2002) 604;
33. V. Cirigliano, G. Paffuti, Commun. Math. Phys. 200 (1999) 381;
34. J. Froelich, P.A. Marchetti, Nucl. Phys. Proc. Suppl. 106 (2002) 47
35. L. Del Debbio, M. D’Elia, A. Di Giacomo, B. Lucini, G. Paffuti , in preparation