MECHANISMS OF COLOUR CONFINEMENT*

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A critical survey is made of the mechanism of color confinement by "dual superconductivity" of QCD vacuum. Tests by Monte Carlo simulations on a lattice are reviewed.

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1. Introduction

QCD is believed to be the theory of strong interactions. Quarks and gluons, however, are not observed in nature as free particles. This fact leads to the statement that colour is confined, which is a way of saying that asymptotic states are only made of colorless particles.

The obvious question is if colour confinement is built in QCD. Perturbative expansion of QCD is highly infrared singular: no picture of physics at large distances emerges from it. Lattice, being the only known formulation of the theory which does not rely on perturbative expansion, is the right tool to investigate this problem. Evidence that QCD does indeed confine colour was already produced by the early numerical simulations [1]. The Wilson loop of size $Ta \times Ra$ (a is the lattice spacing) behaves at large values of R and T as the exponential of the area.

$$W(R,T) \underset{R,T \gg a}{\simeq} \exp[-\sigma a^2 TR]\,, \tag{1.1}$$

$W(R,T)$ is the parallel transport along the rectangular path C of size $Ta \times Ra$

$$W(R,T) = \exp\left(i \int_C A_\mu \, dx^\mu\right) . \tag{1.2}$$

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It can be proved by general arguments that in the limit of infinite mass of the quarks

$$W(R,T) \sim \exp\big(-V(Ra)Ta\big) \tag{1.3}$$

with $V(r)$ the interaction potential between Q and $\bar{Q}$. Eq. (1.1), (1.3) imply

$$V(r) \underset{r\to\infty}{\simeq} \sigma r \tag{1.4}$$

which means quark confinement.

It would be interesting to understand the mechanism of confinement, *i.e.* if it has analogies with some other known physical phenomenon. There are indications that the mechanism could be "dual superconductivity of the second kind"[2, 3]. In what follows I will discuss

(i) what this evidence consists of
(ii) how the mechanism works
(iii) how the mechanism can be tested on lattice.

2. Strings in hadron physics

The dual resonance models [4], which were based on the so called s - t duality [5] led to the idea that the spectrum of hadrons described by Regge trajectories and their daughters, is the spectrum of a string. The action of the Veneziano model is the area of the world sheet of a string [6]

$$S = \sigma \int d\rho\, d\tau \sqrt{\sum_i \left[\left(\frac{\partial X_i}{\partial \rho}\right)^2 + \left(\frac{\partial X_i}{\partial \tau}\right)^2\right]}. \tag{2.1}$$

Quantizing this action gives parallel straight Regge trajectories of particles plus ghosts. Ghosts disappear in 26 dimensions or in 10 dimensions if fermions are included.

The main result established by the intensive study of dual models was that, whatever the fundamental theory of strong interactions would be, its effective hamiltonian at the scale of 1fm should be that of a string.

Fifteen years later strings became again fashionable, in the frame of a kind of opposite philosophy [7]. String theory is the fundamental theory at very short distances, (Planck distance). (Planck mass)^2 plays the role of σ and the standard model is seen as the effective action at the Fermi scale. Also the language has become highbrow, involving concepts of differential geometry.

Nielsen and Olesen [8], inspired by dual models, put forward the idea that strings could be produced by the mechanism which generates Abrikosov flux tubes in a superconductor of 2nd kind [9, 10]. Nambu developed further the idea, studying the motion of open strings with quarks at their ends [11]. After the advent of QCD the idea of string was revived by G. 't Hooft [2] and S. Mandelstam [3]. QCD vacuum acts as a superconductor of second kind which confines the chromoelectric field in Abrikosov like flux tubes. Much work was also devoted to the $1/N_c$ expansion of QCD [12] in which planar diagrams correspond to free strings, and unitarity corrections are classified in terms of topology of the graphs.

3. Flux tubes from lattice

If an ensemble of equilibrium configurations of QCD on a lattice is cooled by a local updating procedure, short distance fluctuations at the scale of lattice spacing rapidly disappear, while correlations at a distance $\xi \gg 1$ survive a much larger number of cooling steps t_c (t_c stays for computer time), with

$$t_c(\xi) \sim \xi^2 . \tag{3.1}$$

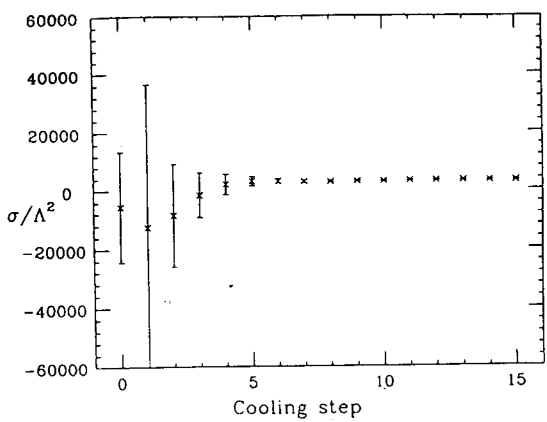


Fig. 1. The SU(2) string tension σ as extracted from $\chi(7,7)$ along cooling, ($\beta = 2.5$). Statistical errors are strongly reduced as short range fluctuations are eliminated. Data correspond to a sample of 50 configurations [13].

The observed behaviour along cooling [13, 14] of the expectation value of the Creutz ratios is shown in Fig. 1. Creutz ratios $\chi(R,T)$ are defined in terms of Wilson loops by the equation

$$\chi(R,T) = -\log \frac{W(R+1,T+1)\,W(R,T)}{W(R,T+1)\,W(R+1,T)} . \tag{3.2}$$

If Eq. (1.1) is satisfied

$$\chi(R,T) = \sigma a^2 . \tag{3.3}$$

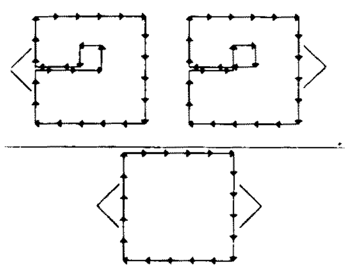
Since a is known as a function of β at large β's, by renormalization group arguments

$$a = \frac{1}{\Lambda} f(\beta)\,, \tag{3.4}$$

$$f(\beta) \simeq \left(\frac{\beta}{2Nb_0}\right)^{b_1^2/2b_0} \exp\left(-\frac{\beta}{2Nb_0}\right) \tag{3.5}$$

$$b_0 = \tfrac{11}{3}\frac{1}{4\pi^2}(N_c - 2N_f)\,, \quad b_1 = \frac{1}{64\pi^4}\left(\tfrac{34}{3}N_c^2 - 2N_f\left(\tfrac{5}{3}N_c + \frac{N_c^2-1}{2N_c}\right)\right)\,,$$

σ/Λ^2 can be extracted from the values of χ produced by numerical simulations. These values are displayed in Fig. 1. Fluctuations due to short range quantum excitations are drastically reduced after a few cooling steps, while the string tension, which is related to long range corrections, survives with small errors and with a value consistent with what is measured without cooling, on much larger samples of configurations [15]. One can then explore [16] the field configurations produced by the propagating $Q\,\bar{Q}$ pair, polished of uninteresting local quantum fluctuations. This is done by measuring the correlations



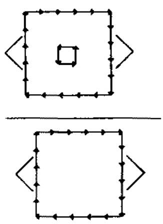
$$\underset{a\to 0}{\simeq} ag\left\{\langle F_{\mu\nu}\rangle_{\bar{Q}Q} - \langle F_{\mu\nu}\rangle_{\mathrm{vac}}\right\} \tag{3.6}$$

The plaquette $\Pi_{\mu\nu}$ measures, (as $a \to 0$), the field strength $F_{\mu\nu}$: the point in which $F_{\mu\nu}$ is measured, as well as the orientations (which correspond to different components of $\vec{E}$ and $\vec{B}$ fields) can be varied, and a map of the field can be constructed. The lines joining the plaquette to the Wilson loop in (3.6) are needed to make the definition of the field strength gauge invariant: in continuum language they correspond to a parallel transport. Of course for an abelian theory they are trivial.

The result of this exploration can be synthetized as follows. At fixed time the component of the chromoelectric field parallel to the line joining the position of $\bar{Q}$ and Q is constant along the line, is maximum at zero transverse distance, and decreases exponentially as the transverse distance is increased

$$E_{\parallel} = E_{\parallel}(0)\exp(-x_T/\lambda)\,, \qquad \lambda \simeq 0.3\,\text{fm}. \tag{3.7}$$

Other components of the electric field and the magnetic field are much smaller. The configuration is shown schematically in Fig. 2. An alternative way of exploring the configuration is by measuring the correlation



$$- \langle \square \rangle \underset{a\to 0}{\simeq} a^2 g^2 \left\{ \langle F^2 \rangle_{\bar{Q}Q} - \langle F^2 \rangle_{\text{vac}} \right\} \tag{3.8}$$

This correlation is smaller numerically than the correlation (3.6) since it is proportional to $(a^2 F_{\mu\nu})^2$, and $a^2 F_{\mu\nu}$ is a small number ($\sim 10^{-1}$) in the scaling region. It has also been measured without cooling [17] and indicates configurations of the same shape as the correlations of Fig. 2. The net result of these investigations is that chromoelectric flux tubes do exist in the region of space between $\bar{Q}\ Q$ pairs.

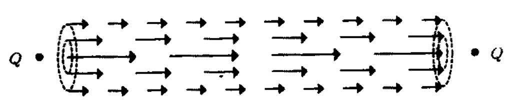


Fig. 2. The chromoelectric field in the region of space between heavy $Q\ \bar{Q}$ pair. Transverse size of flux tube is $\lambda \sim .5\,\text{fm}$ [16].

It would be nice to study not only the lowest configuration of such strings, but also their excitations. Some attempts do exist in the literature in this direction, [18], even if the distances which can be explored in numerical simulations are only few lattice spacings, and therefore the excitations severely distorted with respect to continuum.

4. Superconductivity. Meissner effect

Let us consider the Higgs model, which is a relativistic version of the Landau Ginzburg equation.

$$\mathcal{L} = (D_\mu\varphi)^\dagger (D_\mu\varphi) - V(\varphi) - \tfrac{1}{4}F_{\mu\nu}F^{\mu\nu}\,. \tag{4.1}$$

In this model the photon field is coupled minimally to a scalar charged field φ:

$$D_\mu = \partial_\mu + ieA_\mu \tag{4.2}$$

and

$$V(\varphi) = -\frac{\mu^2}{2}\varphi^*\varphi + \frac{\lambda}{4}(\varphi^*\varphi)^2 . \tag{4.3}$$

If $\mu^2 > 0$, since λ must be positive to have a stable vacuum, the shape of the potential $V(\varphi)$ as a function of $\varphi_1 = \mathrm{Re}\varphi$ $\varphi_2 = \mathrm{Im}\varphi$ is shown in Fig. 3. The theory (4.1) is invariant under the gauge transformation

$$\begin{aligned} \varphi &\to \varphi\, e^{-ie\alpha(x)} , \\ A_\mu &\to A_\mu - \partial_\mu\alpha(x) . \end{aligned} \tag{4.4}$$

The ground state corresponds to a minimum of $V(\varphi)$, (4.3). The solutions of the equation $\partial V(\varphi)/\partial\varphi = 0$, are $\varphi = 0$, which corresponds to the maximum in Fig. 3 and

$$\bar{\varphi} = \sqrt{\frac{\mu^2}{\lambda}}\, e^{i\delta} \tag{4.5}$$

with arbitrary δ. Any choice of δ makes the ground state U(1) non invariant. A situation in which a theory is invariant under a group of transformations, but its ground state is not, is usually called a spontaneous breaking of the symmetry. The above Higgs model spontaneously breaks U(1) invariance. The physical consequence of this breaking is superconductivity. If we put

$$\varphi = \psi\, e^{i\chi e} , \tag{4.6}$$

a gauge transformation (4.4) acts as follows

$$\psi \to \psi , \qquad \chi \to \chi + \alpha . \tag{4.7}$$

In terms of the fields Eq. (4.6)

$$D_\mu\varphi = e^{i\chi}\left[\partial_\mu + ie(A_\mu - \partial_\mu\chi)\right]\psi . \tag{4.8}$$

Putting

$$\tilde{A}_\mu = A_\mu - \partial_\mu\chi , \tag{4.9}$$

$$\mathcal{L} = \partial_\mu\psi\partial^\mu\psi + V(\psi) + e^2\psi^2\tilde{A}_\mu\tilde{A}^\mu - \tfrac{1}{4}\tilde{F}_{\mu\nu}\tilde{F}^{\mu\nu} . \tag{4.10}$$

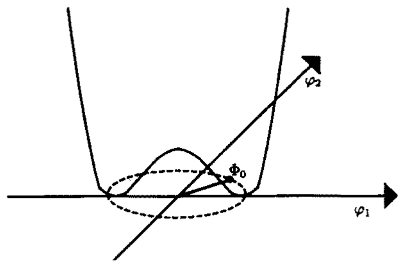


Fig. 3. Potential $V(\varphi)$ of the Higgs model.

All the quantities appearing now in Eq. (4.10), including $\tilde{A}_\mu$ are gauge invariant. Putting

$$\psi = \sqrt{\frac{\mu^2}{\lambda}} + \delta\,, \tag{4.11}$$

the equation of motion for $\tilde{A}_\mu$ becomes, at the lowest order in the interaction with the field δ,

$$\partial^\mu F_{\mu\nu} + \frac{e^2\mu^2}{\lambda}\tilde{A}_\nu = 0 \tag{4.12}$$

which is the equation of a massive vector field. Eq. (4.12) is equivalent to

$$(\Box + m^2)\,\tilde{A}_\mu = 0\,, \qquad \partial^\mu \tilde{A}_\mu = 0\,. \tag{4.13}$$

The first of these equations describes Meissner effect: the field has a finite penetration depth into the medium. Deep inside the vacuum, *i.e.* far from currents $\tilde{A}_\mu = 0$. Integrating on a path around a flux tube

$$\oint \tilde{A}_\mu\, dx^\mu = 0\,, \tag{4.14}$$

or, by Eq. (4.9)

$$\oint A_\mu\, dx^\mu = \Delta\chi\,. \tag{4.15}$$

The difference in the value of χ, $\Delta\chi$, after a winding can only be such that

$$e\Delta\chi = 2\pi\,, \tag{4.16}$$

by Eq. (4.16) and this means flux quantization.

In a real superconductor

$$\varphi \sim \varepsilon_{ij}\psi^i\psi^j\,, \tag{4.17}$$

is the field describing a Cooper pair, which is a pair of electrons in singlet spin state. As in our relativistic model two lengths appear in a real superconductor: a correlation length of the Higgs field, ℓ_1, (μ^{-1} in our model) and the inverse mass of the photon $\ell_2 = \sqrt{\lambda}/\mu e$. If $\ell_1 \gg \ell_2$ the superconductor is first kind: when the magnetic field becomes higher than the critical value the field fills the sample and Meissner effect disappears. If $\ell_2 \gg \ell_1$ the superconductor is 2-nd kind. As the magnetic field grows above some critical value, as shown first by Abrikosov [9], it becomes energetically advantageous that the field penetrates through flux tubes of transverse section ℓ_1, whose density becomes larger and larger as the magnetic field increases, until they fill the sample and superconductivity disappears.

The question we will address is if the flux tubes observed in QCD are in fact Abrikosov flux tubes of some superconductor. Of course the role of electric and magnetic field has to be exchanged with respect to an ordinary superconductor, since flux tubes, as shown in Sec. 3 are of chromoelectric field: this is the meaning of the word "DUAL".

What superconductor is then QCD vacuum, or what U(1) invariance is spontaneously broken? In analogy with Cooper pairs, monopoles, or more generally operators with non trivial monopole number should condense in the ground state. An additional complication appears, as compared to usual superconductors: QCD vacuum is invariant under charge conjugation, which is an exact symmetry of strong interactions. Moreover the symmetry which must break spontaneously to produce superconductivity cannot correspond to a subgroup of the colour group, which is also known to be an exact symmetry.

Monopoles do exist in gauge theories [19, 20]. In a gauge theory with gauge group SU(2), and a scalar field $\vec{\varphi}$ in its adjoined representation [21]

$$\mathcal{L} = -\tfrac{1}{4}\vec{G}_{\mu\nu}\vec{G}^{\mu\nu} + D_\mu\vec{\varphi}D^\mu\vec{\varphi} - V(\vec{\varphi}^{\,2}) \tag{5.1}$$

with

$$V(\vec{\varphi}^{\,2}) = -\frac{\mu^2}{2}\vec{\varphi}^2 + \frac{\lambda}{4}(\vec{\varphi}^{\,2})^2 . \tag{5.2}$$

Higgs phenomenon takes place . If one looks for solutions with finite energy necessarily

$$|\vec{\varphi}| \underset{|\vec{x}|\to\infty}{\longrightarrow} |\vec{\varphi}_0| \tag{5.3}$$

with $\vec{\varphi}_0$ the field which makes $V(\varphi)$ minimum. Therefore these solutions correspond to mapping of the $2d$ sphere at infinity on SU(2). The homotopy group of such mapping is an integer which counts monopoles. In Ref. [19, 20] an explicit example is found of such solutions known as hedgehog solution.

At large distances the solution behaves as follows

$$\frac{\vec{\varphi}}{|\varphi_0|} \underset{|\vec{x}|\to\infty}{\simeq} \frac{\vec{x}}{|x|} \equiv \vec{n}\,, \tag{5.4}$$

$$\vec{A}_{ij} \underset{|\vec{x}|\to\infty}{\simeq} \varepsilon_{ijk} \frac{x^k}{|\vec{x}^2|}\,. \tag{5.5}$$

In Eq. (5.5) i means space component, j component in the SU(2) space. As in Ref. [19, 22] we define

$$F_{\mu\nu} = \vec{n} \cdot \vec{F}_{\mu\nu} \tag{5.6}$$

and

$$B_\mu = \vec{n} \cdot \vec{A}_\mu\,. \tag{5.7}$$

In the hedgehog solution $F_{\mu\nu}$ is, at large distances, the field strength produced by a Dirac monopole, and is a colour singlet.

Also B_μ is a colour singlet. In general [22]

$$F_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu + \frac{1}{g}(\partial_\mu \vec{n} \wedge \partial_\nu \vec{n}) \cdot \vec{n}\,. \tag{5.8}$$

If we choose a gauge $\vec{n} = (1, 0, 0)$ then the last term in Eq. (5.8) vanishes and

$$F_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu\,, \tag{5.9}$$

$F_{\mu\nu}$ looks like an abelian gauge field, with vector potential B_μ. Such choice of gauge is called an abelian projection.

The field $F_{\mu\nu}$, as well as the charge of the monopole which produces it are singlets of the gauge group and invariant under charge conjugation (Eqs (5.6), (5.7)). The corresponding U(1) symmetry would then be a good candidate for breaking spontaneously, and give the dual superconductivity needed to confine colour.

However, contrary to the above model, QCD has no $\vec{\varphi}$ field. How can monopole be defined and with them an abelian projection? A solution was proposed by 't Hooft [23]. We will consider for the sake of simplicity an SU(2) gauge group: the discussion for SU(3) is completely analogous. The idea is that any operator $\vec{O}$ transforming in the adjoined representation can define an abelian projection and monopoles as follows.

In general we will have

$$O(x) = \vec{\Phi}(x) \cdot \vec{\sigma}\,. \tag{5.10}$$

A gauge transformation which diagonalises $O(x)$, *i.e.* which rotates $\vec{\Phi}(x)$ to $\vec{\Phi}'(x)$ with

$$\vec{\Phi}'(x) = (0, 0, \Phi'_3(x))\,,$$

is singular on the world lines when $\vec{\Phi} = 0$. This singularity corresponds to a mapping of the $2d$ sphere on SU(2), which has the topology of a monopole [24].

Of course the location and the number of such monopoles do depend on the choice of the operator $O(x)$, being related to its zeros. The idea of 't Hooft is that physics, (*i.e.* monopole condensation) should be independent of the choice of $O(x)$.

We will explore two possibilities:

(a) 't Hooft is correct in the sense that, whatever operator $O(x)$ one chooses, the corresponding monopoles give rise to confinement by condensing in the vacuum.

(b) There exists at least one choice of $O(x)$ for which the phenomenon takes place.

6. Results from a lattice. Conclusions

If 't Hooft is right, after abelian projection of any operator $O(x)$, the chromoelectric field $\vec{E}_{\parallel}$ inside the flux tubes observed in the lattice simulations, should be directed along the third axis (diagonal) [27, 28]. The three color components of the parallel field depicted in Fig. 2, can be measured after abelian projection. Contrary to the expectation the orientation of the field is random in SU(2) space.

The test was made for $O(x) = F_{12}(x)$; for $O(x)$ the Polyakov line, defined as the trace of the parallel transport, at a fixed spatial position $\vec{x}$, along the time axis, with negative result.

Equally negative answer is obtained [27] for the so called maximal abelian projection [25], which consists in minimizing by gauge transformations the quantity

$$\mathrm{Tr}\left\{U_\mu(n)\sigma_3 U_\mu^\dagger(n)\sigma_3\right\} .$$

The maximal abelian projection is not of the type suggested by 't Hooft, but can equally well define monopole like singularities of the field strength.

The conclusion is that 't Hooft idea does not work, at least in its wider sense: the answer to question (a) of Sec. 5 is negative.

The possibility (b) is, however, still open.

Many groups [25, 26] have measured on lattice the density of monopoles by looking at the abelian projected magnetic field across the plaquettes at the border of a spatial cube. The flux φ is measured as $\int A_\mu dx^\mu$ along the edges, modulo 2π

$$\varphi = 2\pi n + \bar{\varphi} ,$$

n counts Dirac singular strings of monopoles crossing the plaquette. The net number of strings coming out of a cube counts the monopole inside it.

The idea behind these investigations is that a sudden change in the monopole density as a function of temperature could indicate a transition from a phase with monopole condensation (superconductivity) to a normal phase. On one hand, however, such a definition of monopole density proves to be plagued by lattice artifacts [27], on the other hand the density of monopoles is not a good order parameter for signaling such a phase transition, much in the same way as the density of electrons is irrelevant for superconductivity.

To detect dual superconductivity the vacuum expectation value of an operator with nontrivial monopole charge should be found [27] different from zero in the confined phase, and zero in the deconfined phase.

An operator with these properties $O(x)$ can be constructed. The construction is inspired by Ref. [28]

$$O(x) = \exp\left(i\frac{2\pi}{e}\int d^3\vec{z}\,\vec{B}(\vec{x}-\vec{z})\,\vec{\Pi}(\vec{z})\right), \tag{6.11}$$

with $\Pi^i(\vec{z})$ the conjugate momentum to the field $A^i_\mu(x)$

$$\left[\Pi^i(\vec{x},t), A^j(\vec{y},t)\right] = i\,\delta^{ij}\delta^{(3)}(\vec{x}-\vec{y}), \tag{6.12}$$

$\vec{B}(\vec{x})$ is the vector potential of the field produced by a Dirac monopole

$$B_i(\vec{x}) = \varepsilon_{ij3}\frac{x^j}{|\vec{x}\,|^2}, \tag{6.13}$$

$O(x)$ has non trivial monopole number, as can be easily seen in the Schrödinger picture for the fields, where it acts as a translation operator on the field

$$O(\vec{x})|\vec{A}(\vec{x},t)\rangle = |\vec{A}(\vec{x},t) + \frac{2\pi}{e}\vec{B}(\vec{x})\rangle .$$

Work is in progress to measure $\langle O(x)\rangle$ on the lattice after different abelian projections.

In conclusion

(i) Flux tubes do exist in QCD.

(ii) Not all abelian projections are equivalent.

(iii) The picture of QCD vacuum as a dual superconductor is fascinating and physically appealing. There is a way to test it, by looking at the condensation of operators like $O(x)$ Eq. (6.11). Work is in progress on this subject.

REFERENCES

[1] M. Creutz, *Phys. Rev.* **D 21**, 2308 (1980).
[2] G. 't Hooft in *High Energy Physics*, Proceedings of the EPS Int. Conf., Palermo 1975.
[3] S. Mandelstam, *Phys. Rep.* **23 C**, 245 (1976).
[4] G. Veneziano, *Nuovo Cimento* **57 A**, 190 (1968).
[5] C. Schmidt, *Phys. Rev. Lett.* **20**, 689 (1968).
[6] C. Rebbi, *Phys. Rep.* **12 C**, 1 (1974).
[7] See *eg.* M.B. Green, J.H. Schwartz, E. Witten, *Superstring theory*, Cambridge 1987.
[8] H.B. Nielsen, P. Olesen, *Nucl. Phys.* **B 61**, 45 (1973).
[9] A.B. Abrikosov, *JETP* **5**, 1174 (1957).
[10] For a review of superconductivity in the language of high energy physics see S. Weinberg, *Progr. Theor. Phys. Suppl.* **86**, 43 (1986).
[11] Y. Nambu, *Phys. Rev.* **D 10**, 4262 (1974).
[12] G. 't Hooft, *Nucl. Phys.* **B 72**, 481 (1974).
[13] M. Campostrini, A. Di Giacomo, M. Maggiore, H. Panagopoulos, E. Vicari, *Phys. Lett.* **B 225**, 403 (1989).
[14] For a description of the cooling procedure see A. Di Giacomo, these Proceedings p. 227.
[15] C. Michael, M. Teper, *Nucl. Phys.* **305**, 453 (1988).
[16] A. Di Giacomo, M. Maggiore, S.Olejnik, *Phys. Lett.* **B 236**, 199 (1990); *Nucl. Phys.* **B 347**, 441 (1990).
[17] R.W. Haymaker, J. Wosiek, *Phys. Rev. Rapid Comm.* **D36**, 3297 (1987).
[18] M. Caselle, F. Gliozzi, S. Vinti, *Nucl. Phys. Proc. Suppl.* **B 30**, 889 (1983).
[19] G. 't Hooft, *Nucl. Phys.* **B 79**, 276 (1974).
[20] A.M. Polyakov, *JETP Lett.* **20**, 194 (1974).
[21] H. Georgi, S. Glashow, *Phys. Rev. Lett.* **28**, 1494 (1972).
[22] J. Arafune, P.G.O. Freund, G.J. Goebel, *Journ. Math. Phys.* **16**, 433 (1974).
[23] G. 't Hooft, *Nucl. Phys.* **B 190**, 455 (1981).
[24] A.S. Kronfeld, G. Schierholz, U. Wiese, *Nucl. Phys.* **B 293**, 461 (1987).
[25] A.S. Kronfeld, A.L. Laursen, G. Schierholz, U. Wiese, *Phys. Lett.* **B 198**, 516 (1987).
[26] T.L. Ivanenko, A.V. Pochinsky, M.I. Polikarpov, *Phys. Lett.* **B 252**, 631 (1990).
[27] L. Del Debbio, A. Di Giacomo, M. Maggiore, S. Olejnik, *Phys. Lett.* **B 267**, 254 (1991).
[28] J. Greensite, J. Winchester, *Phys. Rev* **D 40**, 4167 (1989).
[29] E.C. Marino, J.S. Ruiz, *Phys. Rev.* **D 39**, 3690 (1989).