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Monopole-antimonopole interaction in Abelian Higgs model.

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Abstract

We consider interaction of a probe monopole-antimonopole pair in the vacuum of the Abelian Higgs model. For simplicity, the mass of the Higgs particle is assumed to be much larger than the mass of the photon (London limit). In case of a massive photon the straightforward application of the Zwanziger formalism to accommodate both magnetic and electric charges is known to result in gauge dependence and infrared instabilities. We argue that the use of the string representation of the Abelian Higgs model allows to ameliorate both difficulties. In particular, we arrive at a well defined expression for the potential energy of the static monopole sources. We argue that the monopole-antimonopole interaction cannot be described by a massive photon exchange with a definite propagator having simple analytical properties.

1. Magnetic monopoles emerge with necessity as degrees of freedom in Abelian projections of gluodynamics [1]. Moreover, there is ample evidence in the lattice simulations for the relevance and even dominance of the magnetic monopoles in the Maximal Abelian projection (for recent reviews see, e.g., [2]) and supporting the idea of the dual superconductor as the confinement mechanism [3]. Hence, there exists the growing interest in field theoretical approach to description of interaction of monopoles (see, e.g., papers [4] and references therein).

In this note we consider interaction of a static monopole-antimonopole pair brought to the vacuum of the Abelian Higgs model. An analogous problem in case of gluodynamics would be interaction of two static quarks with account of the monopole condensation. We will not, however, exploit this analogy and concentrate on the Higgs model itself.

The problem of evaluation of the interaction energy of a monopole pair in superconductor has its own history ((see [5, 6, 7, 8, 9] and references therein). In the field theoretical approach, the crucial element is the Zwanziger formalism [10] which allows to describe consistently photon interacting with magnetic end electric charges. In the original formulation of the formalism the vacuum is assumed to be trivial. If, on the other hand, one tries to incorporate spontaneous symmetry breaking into the Zwanziger formalism, the resulting photon propagator appears to be gauge dependent and involves unphysical singularities. Various ways to deal with these problem based on physical grounds were proposed [5, 7, 8] but the very idea to invoke new principles to derive the propagator does not look satisfactory from the theoretical point of view. The new ingredient which we propose is to use the string formulation of the Abelian Higgs model [11] which is especially simple if one considers the London limit, $m_H \gg m_V$. This approach allows to evaluate the interaction of the static monopoles in an unambiguous way. Interpretation of the result in terms of the photon propagator is not viable, however.

2. Let us start with an overview of the Zwanziger formalism [10], which represents a version of the local field theory of electrically and magnetically charged particles. One introduces two potentials $A_\mu(x)$ and $B_\mu(x)$ which covariantly interact with electric and magnetic currents, respectively. The partition function of the theory reads:

$$\mathcal{Z}_{Zw} = \int DADB\, e^{-S_{Zw}(A,B)+ie(j^e,A)+ig(j^m,B)} \tag{1}$$

where $e(g)$ is the electric (magnetic) charge, $(j,A) = \int d^4x\, j_\mu(x) A_\mu(x) = \int_C A_\mu d\,x_\mu$, and the action $S_{Zw}(A,B)$ is given by:

$$\begin{gathered} S_{Zw}(A,B) = \int d^4x\,\{\tfrac{1}{2}(n\cdot[\partial\wedge A])^2 + \tfrac{1}{2}(n\cdot[\partial\wedge B])^2 + \\ +\tfrac{i}{2}(n\cdot[\partial\wedge A])(n\cdot[\partial\wedge B]^d) - \tfrac{i}{2}(n\cdot[\partial\wedge B])(n\cdot[\partial\wedge A]^d)\}, \end{gathered} \tag{2}$$

where

$$[A\wedge B]_{\mu\nu} = A_\mu B_\nu - A_\nu B_\mu, \quad (n\cdot[A\wedge B])_\mu = n_\nu (A\wedge B)_{\nu\mu},$$
$$(G)^d_{\mu\nu} = \frac{1}{2}\varepsilon_{\mu\nu\lambda\rho} G_{\lambda\rho}$$

The classical equations of motion are:

$$\begin{gathered} (n\partial)^2 A_\mu - n_\mu (n\partial)(\partial A) - \partial_\mu (n\partial)(nA) + n_\mu \partial^2 (nA) - i(n\partial)\varepsilon_{\mu\nu\lambda\rho} n_\lambda \partial_\rho A_\nu = -iej^e_\mu \\ (n\partial)^2 B_\mu - n_\mu (n\partial)(\partial B) - \partial_\mu (n\partial)(nB) + n_\mu \partial^2 (nB) + i(n\partial)\varepsilon_{\mu\nu\lambda\rho} n_\lambda \partial_\rho B_\nu = -igj^m_\mu \end{gathered} \tag{3}$$

Note that both j^m and j^e are conserved as a consequence of the equations of motion, $\partial j^{m,e} = 0$.

Although the theory contains two gauge fields A_μ and B_μ, it still describes one physical photon with two physical degrees of freedom. This follows from a careful treatment of the Hamiltonian dynamics of the system [10, 5].

We are interested in the interaction energy between classical external sources j^m. Although it may be found directly from the equations of motion, this is not the easiest way since (3) is not diagonal in A and B fields. Instead we integrate out all the fields in (1). First of all, we have to fix the gauge freedom which is present in (2). For the gauge-fixing action we choose

$$S_{g.f.} = \int d^4x \left\{ \frac{M_A^2}{2}(nA)^2 + \frac{M_B^2}{2}(nB)^2 \right\} \tag{4}$$

and since there is no ghosts for this gauge the gauge-fixed action of the theory reads:

$$\begin{gathered} S_{Zw} + S_{g.f.} = \tfrac{1}{2}\,(A, \hat{D}^{(A)}A) + \tfrac{1}{2}\,(B, \hat{D}^{(B)}B) - i\,(B, \hat{K}A) - \\ -ie\,(j^e, A)\ - ig\,(j^m, B) \end{gathered} \tag{5}$$

In the momentum space $\hat{D}$ and $\hat{K}_{\mu\nu}$ are defined by:

$$\begin{gathered} \hat{D}^{(A,B)}_{\mu\nu}(k) = \delta_{\mu\nu}(kn)^2 - (kn)(n_\mu k_\nu + n_\nu k_\mu) + (k^2 + M^2_{A,B})n_\mu n_\nu \\ \hat{K}_{\mu\nu}(k) = (kn)(n \wedge k)^d_{\mu\nu} = (kn)\varepsilon_{\mu\nu\lambda\rho}n_\lambda k_\rho \end{gathered} \tag{6}$$

The integration over B field is straightforward yielding the result:

$$\begin{gathered} \mathcal{Z}_{Zw} = \int DA\; e^{-S(A) + ie(j^e, A)} \\ S(A) = \int d^4x \left\{ \tfrac{1}{4}(\partial \wedge A - g(n\partial)^{-1}[n \wedge j^m]^d)^2 + \tfrac{g^2}{2M_B^2}((n\partial)^{-1}(\partial j^m))^2 + \tfrac{M_A^2}{2}(nA)^2 \right\} \end{gathered} \tag{7}$$

Here $(n\partial)^{-1}(n \wedge j^m)^d$ corresponds to the Dirac string which is parallel to the vector n and is attached to the current j^m.

Performing the last integration over A field we obtain the result:

$$\begin{gathered} \int DADB\; e^{-S_{Zw}(A, B) + ie(j^e, A) + ig(j^m, B)} = e^{-S(j^e, j^m)} \\ S(j^e, j^m) = \tfrac{g^2}{2}\,(j^m, \hat{Q}^{(B)}j^m) + \tfrac{e^2}{2}\,(j^e, \hat{Q}^{(A)}j^e) + \tfrac{i}{4}eg\,(j^e, \hat{T}j^m) \end{gathered} \tag{8}$$

Thus, the propagators of the original model (1) in the momentum space are:

$$\begin{gathered} < A_\mu A_\nu > = < B_\mu B_\nu > = \hat{Q}^{(A,B)}_{\mu\nu}(k) = \tfrac{1}{k^2}(\delta_{\mu\nu} + \tfrac{k^2 + M^2_{A,B}}{M^2_{A,B}}\tfrac{k_\mu k_\nu}{(kn)^2} - \tfrac{1}{(kn)}(k_\mu n_\nu + k_\nu n_\mu)) \\ < A_\mu B_\nu > = \hat{T}_{\mu\nu} = \tfrac{1}{k^2}\varepsilon_{\mu\nu\lambda\rho}\tfrac{n_\lambda k_\rho}{(kn)} = \tfrac{1}{k^2(kn)}([n \wedge k]^d)_{\mu\nu}. \end{gathered} \tag{9}$$

Note that the last term in the $S(j^e, j^m)$ apparently depends on the arbitrary chosen vector n. The condition that this dependence is only superficial yields the Dirac quantization condition $eg = 4\pi m$ [10]. There are general proofs that upon imposing this condition the physical results are in fact independent on the choice of the vector n_μ. There is a condition attached to the proofs, namely, that the particle trajectories do not intersect with the Dirac strings. It might worth noting that the direct use of the propagators derived above in the basis of plane waves may not satisfy the last condition and perturbative expressions may depend on n_μ.

3. When the formalism is combined with the spontaneous breaking of the $U(1)$ symmetry apparent inconsistencies arise. Indeed, let us assume that a charged scalar field acquires a nonvanishing vacuum expectation value or, even simpler, a mass term $\frac{m_V^2}{2}A_\mu^2$ is added to the Lagrangian. Then a straightforward diagonalization of the bilinear terms in the Lagrangian results in the following propagators of A- and B-fields (see, e.g., [5]):

$$\begin{aligned} < B_\mu B_\nu > (k) &= \tfrac{1}{k^2+m_V^2}\left(\delta_{\mu\nu} + \tfrac{m_V^2}{(kn)^2}(\delta_{\mu\nu}n^2 - n_\mu n_\nu) + ...\right), \\ < A_\mu A_\nu > (k) &= \tfrac{1}{k^2+m_V^2}\left(\delta_{\mu\nu} + ...\right) \end{aligned} \tag{10}$$

where the dots stand for terms proportional to k_μ and which can eventually be omitted because of the current conservation. If we evaluate the interaction energy of a monopole pair due to the (massive) photon exchange then the n_μ-dependence does not drop off. Also, the double pole in (kn) causes infrared problems.

All these difficulties arise due to the spontaneous symmetry breaking. Indeed, in case of the trivial, or perturbative vacuum the correction similar to (10) is the first radiative correction to $< B_\mu B_\nu >$ propagator due to a loop of charged scalar particles. In that case the algebra which led to (10) remains to a great extent unchanged, with the following substitutions: the factor $(k^2 + m_V^2)^{-1}$ is to be replaced by k^{-2} and the factor $m_V^2/(kn)^2$ is to be substituted by $k^2/(kn)^2$. As a result, the n_μ-dependence of the propagator still persists but the overall factor in front of the n_μ-dependent term is $(kn)^{-2}$. As is readily seen, this term does not vanish only if the Dirac strings attached to each monopole are directed along the same line and overlap. The corresponding energy would be the self energy of the Dirac string and is unphysical.

The problems with the propagator (10) have been discussed in the literature [7, 8] and prescriptions were proposed to fix the problems basing on physical arguments. In particular in Ref. [7] it was proposed to consider the London limit, $m_H \gg m_V$, take principal value when integrating the pole $(nk)^2$ over k, and to direct n_μ along the line connecting the static monopole and antimonopole. The reason for these prescriptions is to imitate the results of the solutions of the classical equations of motion [6]. In this note we are looking for a resolution of the difficulties within the field theoretical approach, without invoking further hypotheses. It is our understanding that the difficulties with evaluating, say, static energy of the monopole-antimonopole pair within the standard formalism are of principal nature since one cannot avoid overlap of the Dirac strings and trajectories of charged particles in case of the vacuum condensation of charges.

4. We propose to utilize the string formulation [11] of the Abelian Higgs model (AHM). In more detail, the partition function of AHM,

$$\begin{aligned} \mathcal{Z}_{AHM} &= \int DAD\Phi D\bar{\Phi}\; e^{-S_{AHM}(A,\Phi,\bar{\Phi})} \\ S_{AHM}(A,\Phi,\bar{\Phi}) &= \int d^4x \,\left\{ \tfrac{1}{4}(\partial\wedge A)^2 + \tfrac{1}{2}|(\partial - ieA)\Phi|^2 + \lambda((|\Phi|^2-\eta^2))^2 \right\} \end{aligned} \tag{11}$$

in the London limit ($\lambda \to \infty$) can be exactly rewritten in terms of the word-sheet coordinates

$\tilde{X}(\sigma)$ of the closed ANO strings [11]:

$$\lim_{\lambda\to\infty} \mathcal{Z}_{AHM} = \int\limits_{\delta\Sigma=0} D\Sigma\, e^{-S(\Sigma)} \tag{12}$$

The action for the ANO strings reads:

$$\begin{gathered} S(\Sigma) = \frac{\pi^2}{e^2}\, m_V^2\, (\Sigma, \hat{K}\Sigma) = \\ = \frac{\pi^2}{e^2} m_V^2 \int d^2\sigma d^2\sigma'\, \varepsilon^{ab}\partial_a\tilde{X}^\mu(\sigma)\partial_b\tilde{X}^\nu(\sigma)\; K(\tilde{X}(\sigma)-\tilde{X}(\sigma'))\; \varepsilon^{a'b'}\partial_{a'}\tilde{X}^\mu(\sigma')\partial_{b'}\tilde{X}^\nu(\sigma') \end{gathered} \tag{13}$$

where the kernel $K(x)$ satisfies the equation $(-\partial^2 + m_V^2)K(x) = \delta(x)$ and $m_V^2 = e^2\eta^2$.

Consider now the static monopole-antimonopole pair in the vacuum of the Abelian Higgs model. In the framework of the Zwanziger formalism this problem corresponds to the expectation value of the Wilson loop:

$$< H(j^m) >= \frac{1}{\mathcal{Z}} \int DADBD\Phi D\bar{\Phi}\; e^{-S_{Zw}(A,B)+\int d^4x \left\{ \frac{1}{2}|(\partial - ieA)\Phi|^2 + \lambda|(|\Phi|^2-\eta^2)|^2 \right\} + ig\,(j^m,B)} \tag{14}$$

The integral over B is the same as in (1), the result[1] is analogous to eq.(7) (see also ref.[12]):

$$\begin{gathered} < H(j^m) >= \frac{1}{\mathcal{Z}_{AHM}} \int DAD\Phi D\bar{\Phi}\; e^{-S_{AHM}(A,\Phi,\bar{\Phi})}\; H(j^m) \\ H(j^m) = e^{\int d^4x \left\{ -\frac{1}{4}(\partial\wedge A + \frac{2\pi}{e}{\Sigma_C}^d)^2 + \frac{1}{4}(\partial\wedge A)^2 \right\}} \end{gathered} \tag{15}$$

where ${\Sigma_C}^d = (n\partial)^{-1}[n\wedge j^m]^d$, $\delta\Sigma_C = j^m$; and we used Dirac quantization condition $eg = 2\pi$.

Thus the expectation value of the Wilson loop $\exp\{ig(j^m, B)\}$ for the gauge field B is reduced to the expectation value of the 't Hooft loop $H(j^m)$ for the gauge field A. The surface spanned on the loop is parallel to the vector n. Now we show that in the string representation $< H(j^m) >$ does not depend on the shape of this surface. Consider for simplicity the London limit, then integrating over the (nonsingular) phase of the Higgs field and over the gauge field A we have:

$$\lim_{\lambda\to\infty} < H(j^m) > = e^{-\frac{2\pi^2}{e^2}(j^m, \hat{K}j^m)} \int\limits_{\delta\Sigma=0} D\Sigma\, e^{-S(\Sigma + \Sigma_c)} \tag{16}$$

where $(j^m, \hat{K}j^m) = \sum_{\mathcal{C},\mathcal{C}'} \int_{\mathcal{C}} dx_\mu \int_{\mathcal{C}'} dx'_\mu\; K(x-x')$ and the string action $S(\Sigma)$ is the same as in (13). By the change of variable $\Sigma \to \Sigma - \Sigma_c$ the integral over closed surfaces ($\delta\Sigma = 0$) is reduced to the integral over the surfaces bounded by the loop C which corresponds to the current j^m:

$$\lim_{\lambda\to\infty} < H(j^m) > = e^{-\frac{2\pi^2}{e^2}(j^m, \hat{K}j^m)} \int\limits_{\delta\Sigma=j^m} D\Sigma\, e^{-S(\Sigma)} \tag{17}$$

[1]We skip for simplicity the gauge fixing terms.

Thus the dependence on the shape of the Dirac string disappears and $< H(j^m) >$ depends only on the loop C.

5. Now we estimate the static monopole antimonopole potential in the London limit. For the monopole-antimonopole pair located at the distance R one has:

$$j^m_\mu(x) = \delta_{\mu,0}[\delta(\vec{x} - \vec{R}/2) - \delta(\vec{x} + \vec{R}/2)] \tag{18}$$

and the monopole-antimonopole potential is calculated as

$$V(R) = -\frac{1}{T} ln < H(j^m) > \tag{19}$$

As for the first term $\frac{2\pi^2}{e^2}\ (j^m, \hat{K} j^m)$ in (17) it is easy to find that for the contour (18) it gives:

$$\frac{2\pi^2}{e^2}(j^m, \hat{K} j^m) = \ \text{(self energy)}\ - \int dt\ \frac{\pi}{e^2}\ \frac{e^{-m_V R}}{R} \tag{20}$$

which results in the Yukawa-type contribution to the potential $V(R)$:

$$V(R) = V_1(R) + V_2(R), \qquad V_1(R) = -\frac{\pi}{e^2}\ \frac{e^{-m_V R}}{R} \tag{21}$$

As for the second term $V_2(R)$ it is much more involved. Even at the classical level one has to find the surface bounded by the contour C for which the action $S(\Sigma)$ is minimal. It is difficult to find such a surface in general case but for the action (13) we expect that it should be the surface of minimal area. For the loop defined by (18) the minimal surface is the flat surface parameterized as follows:

$$\begin{array}{ll} \tilde{X}^0 = t & t \in (-\infty; +\infty) \\ \tilde{X}^i = \frac{1}{2} R^i \sigma & \sigma \in (-1; +1) \end{array} \tag{22}$$

The calculation of the action $S(\Sigma^{min.}_C)$ is straightforward. Since for $\Sigma^{min.}_C$:

$$\varepsilon^{ab} \partial_a \tilde{X}^\mu(\sigma) \partial_b \tilde{X}^\nu(\sigma) = (\delta^{\mu,0}\delta^{\nu,i} - \delta^{\nu,0}\delta^{\mu,i}) \frac{R^i}{2} \tag{23}$$

one has:

$$\begin{aligned} S(\Sigma^{min.}_C) &= \tfrac{1}{2}\left[\tfrac{\pi R m_V}{e}\right]^2 \int dt dt' d\sigma d\sigma'\ \frac{d^4k}{(2\pi)^4}\ \frac{1}{k^2 + m_V^2}\ e^{ik(\tilde{X}(\sigma) - \tilde{X}(\sigma'))} = \\ &= \tfrac{1}{2}\left[\tfrac{4\pi R m_V}{e}\right]^2 \int dt\ \frac{d^3k}{(2\pi)^3}\ \frac{\sin^2(\vec{k}\vec{R}/2)}{\vec{k}^2 + m_V^2}\ \frac{1}{(\vec{k}\vec{R})^2} \end{aligned} \tag{24}$$

Collecting all the above we have for the static monopole-antimonopole potential:

$$V(R) = -\frac{\pi}{e^2}\frac{e^{-m_V R}}{R} + \frac{1}{2}\left[\frac{4\pi R m_V}{e}\right]^2 \int \frac{d^3k}{(2\pi)^3}\ \frac{\sin^2(\vec{k}\vec{R}/2)}{\vec{k}^2 + m_V^2}\ \frac{1}{(\vec{k}\vec{R})^2} \tag{25}$$

Choosing the Higgs mass M_H as the UV cut-off we finally obtain:

$$V(R) = -\frac{\pi}{e^2}\,\frac{e^{-m_V R}}{R} + \frac{\pi m_V^2}{2e^2}\,\left[\,R\ln\frac{M_H^2}{m_V^2} - \frac{2}{m_V} + \int_0^\infty dx\,\frac{e^{-R\sqrt{x+m_V^2}}}{[x+m_V^2]^{3/2}}\,\right]$$

$$= \frac{\pi m_V}{2e^2}\,\left\{-2\,\frac{e^{-m_V R}}{m_V R} + m_V R\,\ln\left[M_H^2/m_V^2\right] - 2\left[1-e^{-m_V R}\right] + 2m_V R\,\mathrm{Ei}[m_V R]\right\} \tag{26}$$

$$\mathrm{Ei}[x] = -\int_x^\infty \frac{e^{-t}}{t}dt = \mathbf{C} + \ln[x] + \sum_{k=1}^{\infty}\frac{(-x)^k}{kk!}$$

which completes the evaluation of the potential energy of the static monopole pair in the approximation considered. Note that the analogous potential was obtained in Ref. [7], but the additional regularization was used in this paper. Moreover our approach is valid beyond the London limit.

6. Thus, we see that the use of the string formulation does allow to circumvent the difficulties with the photon propagator spelled in detail above. Namely, there is no infrared divergence since the expression for the potential contains the combination $\sin^2(\mathbf{k}\cdot\mathbf{R})/(\mathbf{k}\cdot\mathbf{R})^2$ which is finite if $\mathbf{k}\cdot\mathbf{R}\to 0$. Moreover, the arbitrary vector n_μ of the Zwanziger formalism is fixed to be directed along the line connecting the monopoles. This fixation is determined by the minimality condition of the area of the surface bounded by the world trajectory of the monopole current.

At first sight we could interpret our results in terms of the photon propagator as well. Indeed, we could define the static propagator $D(\mathbf{k}^2)$as:

$$D(\mathbf{k}^2) \;\equiv\; \int\frac{d^3\mathbf{R}}{(2\pi)^3}V(R)exp(i\mathbf{k}\mathbf{R}). \tag{27}$$

We would obtain an expression for $D(\mathbf{k}^2)$ without difficulty. Since we are working in the Euclidean space, the natural guess would be that in the relativistic version of the propagator we should substitute $\mathbf{k}^2$ by k^2 and the whole propagator is $D(k^2)\delta_{\mu\nu} + \frac{k_\mu k_\nu}{k^2}D_1(k^2)$, the last term proportional to $k_\mu k_\nu$ is unimportant due to the current conservation.

The point is that such a construction would not be valid. Indeed, it would imply existence of the potential energy of monopole-monopole interaction proportional to the distance R at large R and of opposite sign as compared to the case of the monopole-antimonopole considered above. Such an interaction is devoid of any physical meaning however. Hence there is no propagator with the standard properties. The reason is that the static interaction $V(R)$ considered above accounts for the effect of the ANO string which is a classical solution to the equations of motion. The classical solutions, on the other hand, do not obey the standard crossing properties inherent to the standard Feynman propagators.

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