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## Self-interaction effects on screening in three-dimensional QED

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**Abstract.** We have shown that self-interaction effects in massive quantum electrodynamics can lead to the formation of bound states of quark–antiquark pairs. A current–current fermion coupling term is introduced, which induces a well in the potential energy profile. Explicit expressions of the effective potential and renormalized parameters are provided.

In a recent letter [1], Abdalla and Banerjee have discussed the confinement and screening problems in three-dimensional QED. They have studied the inter-‘quark’ potential between two static test charges in a theory of dynamical fermions of mass  $m$  coupled to electromagnetism. Their results indicate that for small separation, the quantum potential tends to the classical logarithmic Coulomb potential. However, for large distance the potential tends to zero. This shows quite conclusively the confining and screening nature of the potential. These results also corroborate the two-dimensional results [2] nicely.

This paper is aimed at studying the stability of the above scenario, in three dimensions, in the presence of self-interaction among the fermions. In particular, we have chosen the well studied current–current Thirring interaction. When the Thirring coupling  $g$  is *positive*, the new model shows a marked departure of a qualitative nature from [1] in the short distance regime. In the potential profile, there appears more structure, in the form of a *well*, indicating a strong repulsion below some critical distance. This might lead to stable bound states of the quark–antiquark pair. The large distance behaviour shows the expected screening. For negative  $g$ , nothing of the above dramatic nature occurs, albeit the potential decreases more sharply for short distance.

We formulate the problem along the lines of [1]. The fermion modes in the gauged Thirring model are integrated out to incorporate quantum (fermion loop) effects in the subsequent classical analysis. This bosonization is done in the large- $m$  approximation. The auxiliary field  $B_\mu$ , introduced to linearize the Thirring term, is integrated next, resulting in a generalized Maxwell–Chern–Simons gauge theory [3, 4]. The theory now contains two independent parameters,  $m$  and  $g$ . In order to gain further insight, we expand the results in powers of  $g$  and keep terms up to  $O(1/m, g, g^2, g/m)$ . Surprisingly, the terms linear in  $g$  do not alter the results very much whereas the effect of the higher-order corrections is substantial, as mentioned earlier. This is our main result.

The ideas of screening and confinement play a central role in gauge theory. The computational hurdles in four dimensions compel us to study the lower-dimensional models.

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But one has to extract results which are not artefacts of low dimensionality and which can be carried on to the real world. Previously, in two-dimensional QED, [2] results were obtained indicating screening and confinement for massless and massive fermions, respectively. QCD was studied by [5], where apart from the dynamical fermion mass, the representations of the dynamical fermions and test charges became important. The problems regarding  $\theta$ -vacuum, screening, confinement and chiral condensate in two- and three-dimensional QCD were discussed in [6].

The parent model is

$$L_F = \bar{\psi} i \gamma^\mu (\partial_\mu - ie A_\mu) \psi - m \bar{\psi} \psi + \frac{g}{2} |\bar{\psi} \gamma^\mu \psi|^2 - \frac{pe^2}{4} |A_{\mu\nu}|^2 + \frac{qe^2}{2} \epsilon_{\mu\nu\lambda} A^\mu A^{\nu\lambda} + J_\mu A^\mu. \quad (1)$$

Here  $A_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ ,  $J_\mu$  is an external conserved current and conventionally one takes  $p = 1/e^2$ ,  $q = \mu/(2e^2)$ . We have considered them arbitrary to keep track of them. The above model is rewritten with the auxiliary field  $B_\mu$  as

$$L_F = \bar{\psi} i \gamma^\mu (\partial_\mu - ie A_\mu - i B_\mu) \psi - \frac{1}{2g} |B_\mu|^2 - m \bar{\psi} \psi - \frac{pe^2}{4} |A_{\mu\nu}|^2 + \frac{qe^2}{2} \epsilon_{\mu\nu\lambda} A^\mu A^{\nu\lambda} + J_\mu A^\mu. \quad (2)$$

The bosonized Lagrangian to  $O(1/m)$  is

$$L_B = -\frac{a}{4} |B_{\mu\nu}|^2 + \frac{\alpha}{2} \epsilon_{\mu\nu\lambda} B^\mu B^{\nu\lambda} - \frac{1}{2g} |B_\mu|^2 - \frac{(a+p)}{4} e^2 |A_{\mu\nu}|^2 + \frac{(\alpha+q)}{2} e^2 \epsilon_{\mu\nu\lambda} A^\mu A^{\nu\lambda} - \frac{ae}{2} A_{\mu\nu} B^{\mu\nu} + e\alpha \epsilon_{\mu\nu\lambda} B^\mu A^{\nu\lambda} + J_\mu A^\mu \quad (3)$$

where  $\alpha = -1/(8\pi)$  and  $a = -1/(6\pi m)$ . Hence we define the effective action as

$$Z(A_\mu, B_\mu) = \int \mathcal{D}A_\mu \mathcal{D}B_\mu \delta(\partial^\mu A_\mu) \exp\left(-i \int d^3x L_B\right). \quad (4)$$

Note that a Lorentz gauge condition is introduced in the measure since the theory has a  $U(1)$  gauge invariance in the  $A_\mu$  field.

The above theory is quadratic in  $B_\mu$  and after a formal integration of it, we arrive at the following effective action consisting of the gauge field  $A_\mu$  and the external current:

$$Z(A_\mu) = \int \mathcal{D}A_\mu \delta(\partial_\mu A^\mu) \exp(-i/4) \int d^3x \left[ 2e^2 A_\mu \frac{4\alpha^2 g - (a+p) + a^2 g \partial^2}{(ag\partial^2 - 1)^2 + 4\alpha^2 g^2 \partial^2} \times (\partial^2 g^{\mu\nu} - \partial^\mu \partial^\nu) A_\nu + e^2 A_\mu \frac{-4(\alpha+q) + 8g(1-\alpha)\partial^2(2\alpha^2 g - a + a^2 g \partial^2)}{(ag\partial^2 - 1)^2 + 4\alpha^2 g^2 \partial^2} \times \epsilon^{\mu\nu\lambda} \partial_\nu A_\lambda - 4J_\mu A^\mu \right]. \quad (5)$$

This is a generalized Maxwell–Chern–Simons-type theory [3,4,7]. For  $g = q = 0$ , this action reduces to the one in [1]. From the  $A_\mu$  equation it follows that  $\partial_\mu A^\mu = 0$ . Let us from now on work with the truncated version of this model keeping only terms of  $O(a, g, g^2, ag)$ . The  $A_\mu$  equation of motion in the Lorentz gauge is

$$\begin{aligned} P \epsilon_{\mu\nu\lambda} \partial^\nu A^\lambda + Q \partial^2 A_\mu + J_\mu &= 0 \\ P &= 2e^2 [(\alpha+q) - 8(2\alpha^2 g^2 - ag)(1+q)\partial^2] \\ Q &= -e^2 [(4\alpha^2 g - a - p) + 2p(2\alpha^2 g^2 - ag)\partial^2]. \end{aligned} \quad (6)$$

In order to solve this equation for  $A_\mu$  in the static limit, let us define the *curl* of  $A_\mu$  as in [1],

$$\mathcal{A}_\mu = -\epsilon_{\mu\nu\lambda} \partial^\nu A^\lambda.$$

This allows us to rewrite (6) as

$$P \epsilon^{\alpha\beta\mu} \partial_\beta \mathcal{A}_\mu + Q \partial^2 \mathcal{A}^\alpha = \epsilon^{\alpha\beta\mu} \partial_\beta J_\mu.$$

The operator acting on  $A_\mu$  can be formally inverted to yield

$$A_\sigma = \frac{P}{P^2 + Q^2 \partial^2} J_\sigma + \frac{Q}{P^2 + Q^2 \partial^2} \epsilon_{\sigma\mu\nu} \partial^\mu J^\nu$$

which reduces to [1]

$$\left( \partial^2 + \left( \frac{P}{Q} \right)^2 \right) (-\epsilon_{\mu\nu\lambda} \partial^\nu A^\lambda) = \frac{P}{Q^2} J_\mu + \frac{1}{Q} \epsilon_{\mu\nu\lambda} \partial^\nu J^\lambda. \quad (7)$$

For  $\mu = 2$  in the static case, (7) reduces to

$$\left( \partial^2 + \left( \frac{P}{Q} \right)^2 \right) A_0 = \frac{1}{Q} J_0. \quad (8)$$

This equation is inverted to get  $A_0$ . We now proceed exactly as in [1]. The potential energy between two external static charges, (emerging from  $J_0$ ), a distance  $L$  apart, is

$$V(L) = -(L_q - L_0) = -q[A_0(x^1 = -L/2, x^2 = 0) - A_0(x^1 = L/2, x^2 = 0)]. \quad (9)$$

The subscript  $q$  denotes the presence of the external charge  $q$ . It is worthwhile to point out that the quantum corrections to this classical (Coulomb) energy are incorporated at the one-loop level of the underlying fermions since we are using the effective action obtained by first integrating the fermions and subsequently the auxiliary field  $B_\mu$ . Only one-loop corrections are present since the fermion determinant was computed by treating the gauge field  $A_\mu$  as external.

Solving the  $A_\mu$ -equation for  $A_0$  we get

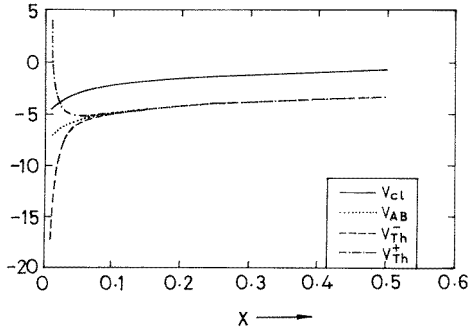
$$A_0(x) = q_1[\Delta(x^1 + L/2, x^2; N) - \Delta(x^1 - L/2, x^2; N)] \\ + q_2 \partial_i \partial_i [\Delta(x^1 + L/2, x^2; N) - \Delta(x^1 - L/2, x^2; N)]. \quad (10)$$

$\Delta(x, m)$  is the two-dimensional Euclidean Feynman propagator given by the modified Bessel function

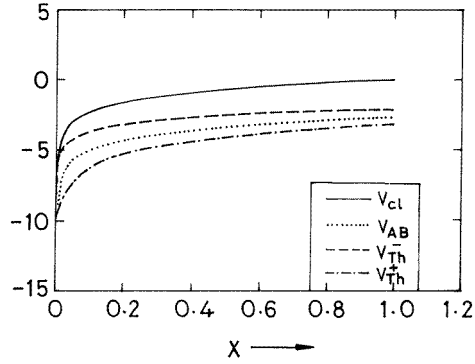
$$\Delta(x, m) = 1/(2\pi) K_0 \left( m \sqrt{(x^1)^2 + (x^2)^2} \right).$$

Here  $q_1$  and  $q_2$  and  $N$  are related to the external charge  $q$  and the fermion mass  $m$ . Once again, for  $g = q = 0$ , we obtain the results of [1]. The interesting term is the latter one in (8) which disappears if *only*  $O(a, g)$  terms are kept. Because of the derivative operator, it radically alters the functional form of  $A_0$ , even though its strength  $q_2$  is much less than  $q_1$ . For small separation  $L$ ,  $K_0(L) \approx \ln(L)$  and the second term  $\approx 1/L^2$ , which competes and finally dominates over the first term. Hence the higher corrections appearing from the self-interaction of the fermions tend to stabilize the charge system. This is the conclusion mentioned at the beginning. Finally, the potential energy function is expressed in terms of dimensionless variables  $M \equiv e^2/m$ ,  $G \equiv g e^2$ ,  $\theta \equiv \mu/e^2$ ,  $X = e^2 L$  as

$$V_{Th} = -\frac{q^2}{\pi} \frac{(1+t)}{1 - M/(6\pi)} K_0 \left[ \frac{(1+u)X}{4\pi(1 - M/(6\pi))} \right] \\ + \frac{q^2}{\pi} \frac{2[G^2/(32\pi^2) + GM/(6\pi)]}{1 - M/(3\pi)} (\partial_X)^2 K_0 \left[ \frac{(1+u)X}{4\pi(1 - M/(6\pi))} \right]. \quad (11)$$



**Figure 1.** The potential functions  $V_{cl}$ ,  $V_{AB}$  and  $V_{Th}^{\pm}$  ( $G = \pm 0.1$ ) are plotted against  $X$ .



**Figure 2.** The potential functions  $V_{cl}$ ,  $V_{AB}$  and  $V_{Th}^{\pm}$  ( $G = \pm 50$ ) are plotted against  $X$ . Only terms linear in  $G$  are considered in  $V_{Th}^{\pm}$ .

This expression is to be compared with  $V_{classical}$  and  $V_{AB}$  in [1]

$$V_{cl} = \frac{q^2}{\pi} \ln(X)$$

$$V_{AB} = -\frac{q^2}{\pi} \frac{1}{1 - M/(6\pi)} K_0 \left[ \frac{X}{4\pi(1 - M/(6\pi))} \right].$$

The correction terms in (11) are

$$u = \frac{4}{1 - M/(2\pi)} \left[ \frac{G}{64\pi^2} + G^2 \left( -\frac{1}{32\pi^3} - \frac{\theta}{128\pi^3} + \frac{\theta}{8\pi^2} - \frac{3\theta^2}{512\pi^3} \right) - MG \left( \frac{1}{96\pi^3} - \frac{1}{6\pi^2} + \frac{\theta}{24\pi^2} - \frac{2\theta}{3\pi} - \frac{\theta^2}{4\pi} \right) \right]$$

$$t = \frac{4}{1 - M/(2\pi)} \left[ \frac{G}{64\pi^2} + G^2 \left( -\frac{1}{1024\pi^4} - \frac{3}{512\pi^3} - \frac{\theta}{64\pi^3} + \frac{\theta}{4\pi^2} + \frac{3\theta^2}{32\pi^2} \right) - GM \left( \frac{1}{64\pi^3} + \frac{1}{3\pi^2} + \frac{\theta}{12\pi^2} - \frac{4\theta}{3\pi} - \frac{\theta^2}{2\pi} \right) \right].$$

This constitutes our main result.

Notice that  $\theta$  does not play any significant role in the present context and so we put  $\theta = 0$ . In figure 1, for  $M = 0.1$ ,  $G = \pm 0.1$ ,  $V_{cl}$ ,  $V_{AB}$  and  $V_{Th}^{\pm}$  are plotted. For this case, the potential well in  $V_{Th}^+$  ( $G = +0.1$ ) is at  $X \approx 2\sqrt{G^2/(32\pi^2)} + GM/(6\pi) \approx 0.046$  and for  $X < 0.003$ ,  $V_{Th}^+$  becomes positive.  $V_{Th}^-$  ( $G = -0.1$ ) shows a sharper descent at short distance. Figure 2 shows only  $O(G)$  corrections to the  $V_{AB}$  result. The fact that  $V_{Th}^{\pm} \approx V_{AB}$  up to  $O(G)$  is quite insensitive for a wide range of values of  $G$  and  $M$ . Here, an unrealistic value  $G = \pm 50$  is chosen just to separate the  $V_{Th}^{\pm}$  and  $V_{AB}$  lines as well as to stress the stability of this formulation up to  $O(M, G)$ . The shift of  $V_{Th}^{\pm}$  with respect to  $V_{AB}$  in the upward or downward direction is dictated by the sign of  $G$ .

Up to  $O(M, G)$  the effective mass of the gauge particle  $A_\mu$  and the renormalized charge  $q^2$  are

$$M_{Th} = \frac{P}{Q} \approx \frac{e^2/(4\pi)}{1 - M/(6\pi)} \left( 1 - \frac{G}{2\pi} \right) = M_{AB} \left( 1 - \frac{G}{2\pi} \right)$$

$$(q^2)_{Th} \approx \frac{q^2}{1 - M/(6\pi)} \left( 1 + \frac{G}{16\pi^2} \right) = (q^2)_{AB} \left( 1 + \frac{G}{16\pi^2} \right).$$

Higher-order momentum dependent terms have been left out.

Let us summarize our results. The potential energy between the external charges at large distance is screened as the gauge particles acquire mass from the Chern–Simons term. Without the Thirring interaction, at short distance, one gets a decreasing negative potential, logarithmic in nature. The positive Thirring coupling term introduces a sort of centrifugal barrier in the effective potential, which leads to the potential well formation. Expressing the centrifugal term as  $(\text{angular momentum})^2/(2ML^2)$  shows clearly that *only* the Thirring term contributes to the correction in the angular momentum of the ‘bound state’. The physical reason is the following. From (2), notice that  $B_\mu/g$  is identified as the fermion current [4]. The *vector nature* of the interaction leads to the derivative term upon integration, which subsequently changes the angular momentum. Estimates of effective mass of the state can be obtained from harmonic oscillator excitations around the well minimum.

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