Chiral Symmetry Breaking in an External Field*

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Abstract

The effects of an external field on the dynamics of chiral symmetry breaking are studied using quenched, ladder QED as our model gauge field theory. It is found that a uniform external magnetic field enables the chiral symmetry to be spontaneously broken at weak gauge couplings, in contrast with the situation when no external field is present. The broken chiral symmetry is restored at high temperatures as well as at high chemical potentials. The nature of the two chiral phase transitions is different: the transition at high temperatures is a continuous one whereas the phase transition at high chemical potentials is discontinuous.

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I would like to report here some recent work [1,2] done in collaboration primarily with D.-S. Lee and Y. J. Ng. The objective is to understand how an external field may affect the dynamics of chiral symmetry breaking in gauge theories. Quantum electrodynamics (in $3 + 1$ dimensions), treated in the quenched, ladder (or planar) approximation, will serve as our model gauge theory. At present our formalism is applicable only for a constant (in space as well as in time) external field and I shall discuss results pertinent to the case of a constant magnetic field.

The problem of dynamical chiral symmetry breaking in a uniform magnetic field in quenched, planar QED has received some attention in recent years [3] - [13]. Following the pioneering works [14] on the study of chiral symmetry breaking in quenched, ladder QED, our approach is based on the Schwinger-Dyson (SD) equation for the fermion self-energy. In this approach, one starts with a Lagrangian with a zero bare fermion mass and looks for nontrivial solutions to the SD equation which signal a dynamically generated fermion mass and a possible spontaneous breakdown of the chiral symmetry.

Let us consider QED with a single charged fermion having a zero bare mass and charge $g$. In the quenched, ladder approximation, the SD equation in the $x$-representation reads

$$M(x, x') = ig^2 \gamma^\mu G_A(x, x') \gamma^\nu D_{\mu\nu}(x - x'),$$

where $M(x, x')$ is the fermion mass operator in the $x$-representation, $M(x, x') = (x|\hat{M}|x')$, $D_{\mu\nu}(x - x')$ is the bare photon propagator, and $G_A(x, x')$ is the fermion propagator in the presence of an external field represented by the vector potential $A_\mu(x)$. We adopt the metric with the signature $g_{\mu\nu} = (-1, 1, 1, 1)$. $G_A$ satisfies the equation,

$$\gamma^\mu \Pi_\mu G_A(x, y) + \int d^4x' M(x, x') G_A(x', y) = \delta(4)(x - y),$$

where $\Pi_\mu \equiv -i\partial_\mu - gA_\mu(x)$. Schwinger [15] was the first to obtain an exact analytical expression for the fermion Green’s function in the presence of a constant electromagnetic field of arbitrary strength. However, we find the alternative representation of $G_A(x, y)$ proposed by Ritus [16] more convenient for our purpose. The essence of this approach is explained below.
Since, in the presence of a constant external field, the fermion asymptotic states are no longer free particle states represented by plane waves, but are described by wavefunctions consistent with the particular external field configuration, namely, eigenfunctions of $\left(\gamma^\mu \Pi_\mu\right)^2$:

$$- \left(\gamma \cdot \Pi\right)^2 \psi_p(x) = p^2 \psi_p(x).$$ (3)

Instead of the usual momentum space, it is more convenient to work in the representation spanned by these eigenfunctions. Another advantage of using this representation is that, for constant external fields, the mass operator is diagonal [16].

If we work in the chiral representation of the Dirac matrices in which $\gamma_5$ and $\Sigma_3 = i\gamma_1\gamma_2$ are both diagonal with eigenvalues $\chi = \pm 1$ and $\sigma = \pm 1$, respectively, the eigenfunctions $\psi_p(x)$ has the general form

$$\psi_p(x) = E_{p\sigma\chi}(x)\omega_{\sigma\chi},$$ (4)

where $\omega_{\sigma\chi}$ are bispinors which are the simultaneous eigenvectors of $\Sigma_3$ and $\gamma_5$. The exact functional form of the $E_{p\sigma\chi}(x)$ will depend on the specific external field configuration.

In the case of a constant magnetic field of strength $H$ pointing in the $z$-direction, the vector potential may be taken to be $A_\mu = (0, 0, Hx_1, 0)$ and one finds that the eigenfunctions $E_{p\sigma\chi}(x)$ do not depend on $\chi$:

$$E_{p\sigma}(x) = Ne^{i(p_0x^0 + p_2x^2 + p_3x^3)}D_n(\rho).$$ (5)

Here $N$ is a normalization factor and $D_n(\rho)$ are the parabolic cylinder functions [17] with argument $\rho \equiv \sqrt{2|gH|}(x_1 - \frac{p_2}{gH})$ and index $n$ which labels the Landau levels:

$$n = n(k, \sigma) \equiv k + \frac{gH\sigma}{2|gH|} - \frac{1}{2}, \quad n = 0, 1, 2, ...$$ (6)

The eigenvalue $p$ now stands for the four quantum numbers $(p_0, p_2, p_3, k)$, where $k$ is the discrete quantum number of the quantized squared transverse momentum:

$$- \left(\gamma \cdot \Pi_\perp\right)^2 \psi_p(x) \equiv - (\gamma^1\Pi_1 + \gamma^2\Pi_2)^2 \psi_p(x)$$

$$= 2|gH|k\psi_p(x).$$ (7)
For a given \( n \), the allowed values for \( k \) are \( k = n, n + 1 \).

The eigenfunction-matrices \( E_p(x) \) defined as

\[
E_p(x) \equiv \sum_\sigma E_{p\sigma}(x) \text{diag}(\delta_{\sigma 1}, \delta_{\sigma -1}, \delta_{\sigma 1}, \delta_{\sigma -1}) \\
\equiv \sum_\sigma E_{p\sigma}(x) \Delta(\sigma)
\]

(8)
satisfy the orthonormality and completeness relations \((E_p \equiv \gamma^0 E_p^{\dagger} \gamma^0)\):

\[
\int d^4x \bar{E}_{p'}(x) E_p(x) = (2\pi)^4 \delta^{(4)}(p - p') \equiv (2\pi)^4 \delta_{kk'} \delta(p_0 - p'_0) \delta(p_2 - p'_2) \delta(p_3 - p'_3),
\]

(9)

\[
\int d^4p E_p(x) \bar{E}_p(y) \equiv \sum_k \int dp_0 dp_2 dp_3 E_p(x) \bar{E}_p(y) = (2\pi)^4 \delta^{(4)}(x - y),
\]

(10)

provided that the normalization constant in Eq.(5) is taken to be \( N(n) = (4\pi |g_H|)^{1/4}/\sqrt{n!} \).

They also satisfy the useful relation \([16]\):

\[
\gamma \cdot \Pi E_p(x) = E_p(x) \gamma \cdot \bar{p},
\]

(11)

where \( \bar{p}_0 = p_0, \bar{p}_1 = 0, \bar{p}_2 = -\text{sgn}(gH) \sqrt{2|gH|k}, \bar{p}_3 = p_3 \). Note that, in terms of the momentum \( \bar{p} \), the system is effectively a \((2+1)\)-dimensional one.

These properties of the \( E_p \)-functions enable us to introduce the \( E_p \)-representation of the fermion Green’s function:

\[
G_A(p, p') \equiv \int d^4x d^4y \bar{E}_{p'}(x) G_A(x, y) E_p(y) = (2\pi)^4 \delta^{(4)}(p - p') \frac{1}{\gamma \cdot \bar{p} + \bar{\Sigma}_A(\bar{p})}
\]

(12)

where \( \bar{\Sigma}_A(\bar{p}) \) represents the eigenvalue matrix of the mass operator:

\[
\int d^4x' M(x, x') E_p(x') = E_p(x) \bar{\Sigma}_A(\bar{p}).
\]

(13)

It is straightforward to verify that the inverse transform,

\[
G_A(x, y) = \int d^4p \frac{d^4p}{(2\pi)^4} E_p(x) \frac{1}{\gamma \cdot \bar{p} + \bar{\Sigma}_A(\bar{p})} E_p(y),
\]

(14)
satisfies Eq.(2). Eqs. (12) and (14) are the generalization of the well-known relations between coordinate space and momentum space Green’s functions, with the plane wave eigenfunctions in the Fourier transform replaced here by the $E_p$-eigenfunctions in order to account for the presence of the external field. Eq.(12) shows explicitly that the fermion propagator is diagonal (in momentum) in the $E_p$-representation. As stated earlier, the mass operator is also diagonal in this representation:

$$M(p, p') = \int d^4x d^4x' E_\mu(x) M(x, x') E_\mu(x') = (2\pi)^4 \delta^{(4)}(p - p') \Sigma_A(\bar{p}).$$  

(15)

Transforming to the $E_p$-representation, the SD equation, Eq.(1), becomes

$$\hat{\Sigma}_A(\bar{p}) \delta_{kk'} = ig^2 \sum_k \left( \sum_\{\sigma\} J_{nn''}(\bar{q}_\perp) J_{nn'''}(\bar{q}_\perp) \cdot \Delta \gamma^\sigma \Delta'' \gamma_{\mu} \Delta', \right) \frac{1}{\gamma \cdot \bar{p}'' + \Sigma_A(\bar{p}'')} \hat{\Delta}''' \gamma_{\mu} \Delta',$$

(16)

where

$$J_{nn'''}(\bar{q}_\perp) \equiv \sum_{m=0}^{\min(n, n'')} \frac{n!n''!}{m!(n - m)!(n'' - m)!} [isgn(gH)\bar{q}_\perp]^{n + n'' - 2m},$$

(17)

$$\hat{q}_\perp^2 \equiv \frac{q_0^2 + q_1^2}{2|gH|}, \quad \varphi \equiv \arctan \left( \frac{q_2}{q_1} \right),$$

(18)

and the momentum $\bar{p}''$ is given by: $\bar{p}_0'' = p_0 - q_0$, $\bar{p}_1'' = 0$, $\bar{p}_2'' = -sgn(gH)\sqrt{2|gH|k''}$, $\bar{p}_3'' = p_3 - q_3$. Here $n' = n(k', \sigma')$, $n'' = n(k'', \sigma'')$, $\bar{n}'' = n(k'', \bar{\sigma}'')$, $\Delta' = \Delta(\sigma')$, $\Delta'' = \Delta(\sigma'')$, $\bar{\Delta}'' = \Delta(\bar{\sigma}'')$, and the summation over $\{\sigma\}$ means summing over $\sigma$, $\sigma'$, $\sigma''$, and $\bar{\sigma}''$. Eq.(16) is valid in the Feynman gauge. The issue of gauge dependence has been addressed recently in Ref. [12] which shows that, within the approximations used, the solution [1,2] to this equation which will be discussed below satisfies the Ward-Takahashi identities.

A general analytic solution to Eq.(16) for $\Sigma_A(\bar{p})$ is not yet available. However, one can obtain an approximate infrared solution by the following simplifications. First we observe that, due to the factor $e^{-\hat{q}_\perp^2}$ in the integrand, only the contributions from small values of $\hat{q}_\perp$
are important. We may therefore truncate the \( J_{nn''} \) series and keep only the terms with the smallest power of \( \hat{q}_\perp \), i.e., \( J_{nn''}(\hat{q}_\perp) \to n! \delta_{nn''} \); and similarly for the \( J_{h''n'} \). This will be referred to as the small \( \hat{q}_\perp \) approximation and is valid only for weak couplings \( (g^2/4\pi \ll 1) \).

Next, we sum over the spin indices and note that the remaining summation over \( k'' \) involves at most three terms: for \( k > 0 \), \( k'' = k, k \pm 1 \). In the limit \( k = 0 = \bar{p}_\perp \), we keep only the dominant \( k'' = 0 \) term. This is known as the lowest Landau level approximation \[3\]. The SD equation is now simplified to

\[
\Sigma_A(\bar{p}_\parallel) \simeq 2g^2 \int \frac{d^4 q}{(2\pi)^4} \frac{e^{-\hat{q}_\perp^2}}{q^2} \frac{\Sigma_A(\bar{p}_\parallel - q_\parallel)}{(\bar{p}_\parallel - q_\parallel)^2 + \Sigma_A^2(\bar{p}_\parallel - q_\parallel)},
\]

where we have made a Wick rotation to Euclidean space: \( p_0 \to ip_4, q_0 \to iq_4 \). Note that the fermion wavefunction renormalization vanishes in the Feynman gauge, hence the self-energy \( \Sigma_A \) has been replaced by the dynamically generated fermion mass \( \Sigma_A \) in Eq.(19). Except for the exponential factor in the integrand, Eq.(19) has the same form as the corresponding SD equation when the external field is absent. The difference is that only the longitudinal momentum is relevant here. This reduction of dimensions from 4 to 2 has been stressed in Ref. \[3\].

Finally, we consider the \( \bar{p}_\parallel = 0 \) limit and approximate the \( \Sigma_A(q_\parallel) \) in the resultant integrand by \( \Sigma_A(0) \equiv m \) to secure the gap equation,

\[
1 \simeq 2g^2 \int \frac{d^4 q}{(2\pi)^4} \frac{e^{-\hat{q}_\perp^2}}{q^2} \frac{1}{q_\parallel^2 + m^2}
\]

\[
\simeq \frac{g^2}{4\pi^2} |gH| \int_0^\infty dq_\perp e^{-\hat{q}_\perp^2} \frac{\ln(2|gH|q_\perp^2/m^2)}{2|gH|q_\perp^2 - m^2}.
\]

The solution to Eq.(20) has the form

\[
m \simeq a \sqrt{|gH|} e^{-2\pi b/g},
\]

where \( a \) and \( b \) are positive constants of order 1. The \( g \)-dependence of \( m \) clearly indicates the nonperturbative nature of this result.
The above solution for the fermion dynamical mass is consistent with that found by Gusynin et al. [3], who studied the Bethe-Salpeter equation for the bound-state Nambu-Goldstone bosons of the spontaneously broken chiral symmetry. The order parameter of this symmetry breakdown is computed to be [2,8]

\[
\langle \bar{\psi} \psi \rangle \simeq - \frac{|gH|}{2\pi^2} m \ln \left( \frac{|gH|}{m^2} \right) 
\simeq - \frac{2ab}{g\pi} |gH|^{3/2} e^{-2\pi b/g}.
\]

Note that, by allowing the fermion field \( \psi \) to carry quantum numbers of internal symmetries, \( \langle \bar{\psi} \psi \rangle \to \langle \bar{\psi}_i \psi_j \rangle \), the formalism developed above may be applied to study the dynamical breaking of internal gauge symmetries (in quenched, ladder approximation) in the presence of an external field.

Currently there is no experimental evidence for the magnetic field induced dynamical chiral/gauge symmetry breaking found in the above solution. However, suggestive hints are available in excitonic systems. In their recent experiment with the coupled AlAs/GaAs quantum wells [18], Butov et al., found evidence of exciton condensation (condensation of electron-hole pairs analogous to the fermion-antifermion pairing in the spontaneously broken chiral vacuum) in the form of a huge broad band noise in the photoluminescence intensity when a sufficiently strong magnetic field is applied, while no exciton condensation occurs in the absence of the external magnetic field.

As an application, the method developed in the study of chiral symmetry breaking in an external magnetic field has been employed in the study of the effects of external magnetic fields on high-\( T_c \) superconductors [19]. The dynamical chiral symmetry breaking solution found above may also be relevant for the chiral phase transition in heavy-ion collisions or in the electroweak phase transition [3] in the early universe when a large primordial magnetic field is expected to be present [20]. In these cases, the effects of temperature and chemical potential need be incorporated.

It is straightforward to generalize the above formalism to include nonzero temperature \( (T \neq 0) \) and nonzero chemical potential \( (\mu \neq 0) \) effects. One finds that the gap equation,
Eq. (20), becomes

$$1 \simeq \frac{g^2}{8\pi^2} |gH| \int_{-\infty}^{\infty} dq_3 \int_{0}^{\infty} dq_\perp \frac{e^{-q_\perp^2}}{Q_1 Q_2} \cdot \left\{ Q_1 \left[ \frac{\coth(Q_2)}{Q_1^2 - (Q_2 + \mu - i\pi T)^2} + \frac{\coth(Q_2)}{Q_1^2 - (Q_2 - \mu + i\pi T)^2} \right] + Q_2 \left[ \frac{\tanh(Q_1 + \mu)}{Q_2^2 - (Q_1 + \mu - i\pi T)^2} + \frac{\tanh(Q_1 - \mu)}{Q_2^2 - (Q_1 - \mu + i\pi T)^2} \right] \right\},$$

(23)

where \( Q_1^2 \equiv q_3^2 + m_{T\mu}^2 \), \( Q_2^2 \equiv q_3^2 + 2|gH|q_\perp^2 \), and \( m_{T\mu} \) is the infrared dynamical fermion mass which depends on both the temperature and the chemical potential. We use units in which the Boltzmann constant equals 1. Aside from the quenched, ladder approximation, and the small \( q_\perp \) and lowest Landau level approximations, Eq. (23) is exact in its dependence on the coupling constant, the magnetic field, the temperature, and the chemical potential.

We have considered separately the \( \mu = 0 \) and \( T = 0 \) limits of Eq. (23). We examine both analytically and numerically the behavior of the dynamical mass as \( T \) (or \( \mu \)) is varied. In the \( (T \neq 0, \mu = 0) \) case, we find that \( m_{T0} \) decreases monotonically as the temperature is raised and eventually vanishes above a critical temperature, indicating that the chiral symmetry is restored at high temperatures. The critical temperature at which this continuous phase transition takes place is estimated to be [2,6,7]

$$T_c \sim 2\sqrt{2} m_{00},$$

(24)

where \( m_{00} \) is the \( (T = 0, \mu = 0) \) solution found in Eq. (21). The order parameter for this phase transition exhibits similar behaviors:

$$\langle \bar{\psi}\psi \rangle_{T0} \sim -\frac{|gH|}{2\pi^2} m_{T0} \ln \left( \frac{|gH|}{m_{T0}^2} \right) \sim |gH|^{3/2} \left( 1 - \frac{T}{T_c} \right)^{1/2} \ln \left( 1 - \frac{T}{T_c} \right)^{1/2}, \text{ as } T \to T_c^-.$$

(25)

The last expression reflects the behavior of \( m_{T0} \) near \( T_c \):

$$m_{T0} \sim |gH|^{1/2} \left( 1 - \frac{T}{T_c} \right)^{1/2}, \text{ as } T \to T_c^-.$$

(26)
In the \((T = 0, \mu \neq 0)\) case, \(m_{0\mu}\) also vanishes as \(\mu\) is increased beyond a critical value, thus restoring the chiral symmetry at high chemical potentials. However, this chiral phase transition is discontinuous \([2]\):

\[
\langle \bar{\psi}\psi \rangle_{0\mu} \simeq - \frac{|gH|}{2\pi^2} m_{0\mu} \ln \left( \frac{|gH|}{m_{0\mu}} \right),
\]

\[
m_{0\mu} > 0, \quad \mu < \mu_c,
\]

\[
= 0, \quad \mu > \mu_c.
\]

(27)

The critical chemical potential is approximately

\[
\mu_c \simeq \frac{m_{00}}{\sqrt{1 - \frac{2I_1}{I_2}}}
\]

(28)

where

\[
I_1 = \int_{-\infty}^{\infty} dq_3 \int_0^\infty d\tilde{q}_1^2 \frac{e^{-\tilde{q}_1^2}}{Q^2 Q_2 (Q + Q_2)} \left( \frac{1}{Q} + \frac{1}{Q + Q_2} \right),
\]

\[
I_2 = \int_{-\infty}^{\infty} dq_3 \int_0^\infty d\tilde{q}_1^2 \frac{e^{-\tilde{q}_1^2}}{Q Q_2 (Q + Q_2)^2},
\]

(29)

and \(Q^2 \equiv q_3^2 + m_{00}^2\). The integrals \(I_1\) and \(I_2\) are both positive and finite.

With these results, we can evaluate whether the dynamics of chiral symmetry breaking in a magnetic field discussed here may be relevant for the electroweak phase transition in the early universe. For instance, the electroweak phase transition took place at a temperature of order 100 GeV. From Eq.(24), this requires a magnetic field of order \(2 \times 10^{41}\) gauss if we take \(a = b = 1\) and \(4\pi/g^2 \simeq 137\). This is much larger than any estimates of the magnetic field strength at the time of the electroweak phase transition \([20]\). We may therefore conclude that the chiral symmetry breaking solution considered here does not play any role in the electroweak phase transition.

On the other hand, as suggested earlier, the formalism described here can be used to study the dynamical breaking of internal gauge symmetries, abelian as well as nonabelian, in the presence of a magnetic field. Within the quenched, planar approximation, we expect
the same generic results as obtained here to be applicable in those situations. We may therefore entertain the possibility that the coupling constant is relatively large, e.g., $4\pi/g^2$ of order $0.1$ (As concrete examples, we may consider electroweak symmetry breaking in technicolor models or chiral symmetry breaking in QCD by colored fermions belonging to large representations of SU(3)). In this case the required magnetic field could be of order $10^{28}$ gauss or less, which makes it an interesting possibility for the study of phase transitions in the early universe.

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