RAL-92-038

Science and Engineering Research Council

Rutherford Appleton Laboratory

Chilton DIDCOT Oxon OX11 0QX

RAL-92-038

Heavy-ion collisions and a new phase of QED

E Papageorgiu

June 1992

RAL-92-038

Heavy-ion collisions and a new phase of QED.

E. Papageorgiu
Rutherford Appleton Laboratory
Chilton, Didcot, Oxon
OX11 0QX, England

Abstract

We argue that the "anomalous" narrow peaks, seen in the e^+e^- spectrum of heavy-ion collisions, at energies near the Coulomb barrier, at GSI, may be due to the fact that the effective coupling of the e^+e^- system to the coherent EM-field of the two nuclei, can be of order one at impact parameters comparable (or smaller) to the e^- Compton wavelength. It is conceivable that the excitation of a large number of virtual and real pairs within a Compton volume could result in momentum correlations for the e^+ and e^-, or even lead to a "new phase" of QED with a different mass scale for the electron. The presence of similar anomalies in the e^+e^- spectrum of relativistic, but exclusively coherent, collisions of heavy nuclei, would support this idea and rule out any interpretation based on spontaneous emission from a supercritical QED vacuum. In addition, the use of e^+e^- interferometry could help in clarifying the crucial point of whether the emitted pairs stem from a new or the well-known phase of QED.

(June 1992)

In the last decade, there have been a number of experiments with heavy ions at bombarding energies close to the Coulomb barrier, by two experimental groups (the EPOS and the Orange Spectrometer group) at the Gesellschaft für Schwerionenforschung (GSI) in Darmstad, which have seen some anomalies in the energy spectrum of the emitted positrons and electrons measured in coincidence [1]. The so-called "anomalous peaks" consist of a set of (up to three) narrow lines with a width $\Gamma_{e^\pm} \simeq 20 - 40$ keV, superimposed on the much broader continuum for the induced e^+e^- pair production ($\Gamma_{e^+e^-} \simeq 500$ keV) and the nuclear background, and lie between 1.6 and 1.8 MeV, with only a slight dependence of their position on the energy, the scattering angle, and the combined charge $Z_u = Z_1 + Z_2$ of the nuclei.

When the first peak was discovered in the spectrum of the $U + Cm$ system, whose combined charge exceeds the critical value $Z_{cr} \simeq 172$, for which the binding energy for an electron in the 1s state equals twice its rest mass, people thought that it was the signal for the spontaneous emission of a positron due to the decay of the normal QED vacuum [2]. However, this is possible only if the nuclei can stick together, forming a giant dinuclear complex, for an unusually long time $\sim 10^{-19}$ sec as compared with the collision time of 10^{-21} sec [3]. Even so, the position of the peaks should depend strongly on the value of the charge, scaling $\sim Z^{20-22}$ [3], a fact which is in conflict with what has been measured. Finally, the "anomaly" became obvious when the same lines were found also in subcritical systems.

Another remarkable characteristic of the data is the emission of the e^+ and e^- from a single system, i.e. not from the individual nuclei, which decays at rest in the c.m.s. and results mostly in a back-to-back emission. This, and the absence of anomalies in the time-reversed process in Bhabha ($e + e- \longrightarrow e + e-$) scattering [4], has ruled out interpretations based on the existence of elementary particles decaying into $e + e-$ pairs, and reinforced the idea of the formation of a collective state due to the presence of the background electromagnetic (EM) field of the nuclei. In particular, using the analogy of the positronium as an e^+e^- bound state, the energy-correlated e^+ and e^-, seen at GSI, have been interpreted as bound states in a "new" non-perturbative phase of QED [5], in which chiral symmetry might be broken [6], thus generating a new mass scale for the electron.

The existence of a new non-perturbative phase of QED has been long postulated as the remedy to the problem of the triviality of perturbative QED, which started, when Landau pointed out, that the renormalized coupling α_r, in a one-loop perturbative evaluation of the photon propagator,

$$\alpha_r = \frac{\alpha}{1 + \frac{\alpha}{3\pi} ln(\frac{\Lambda^2}{m^2})} \tag{1}$$

vanishes when taking the momentum cut-off Λ to infinity [7]. A large number of studies have been performed since then on the lattice [8], and analytically [9], by solving the Schwinger-Dyson equations (SDE) in the ladder, quenched, approximation, indicating the existence of a chiral phase transition when the coupling constant becomes of order one. Because at the critical point, four-fermion operators do not decouple, they've been lately added to the gauge-interaction Lagrangian. The corresponding phase diagram exhibits a critical line and a shift to a weaker gauge coupling as the four-fermion interaction is turned on [10].

However, since in heavy-ion collisions, from sub-Coulomb to ultra-relativistic energies, the value of α remains practically equal to $1/137$, there is at first little hope that the anomalous $e+e-$ anomalous peaks are related to the strong-coupling phase of QED, unless an appropriate four-fermion interaction is induced by the ions background electromagnetic (EM) field, for which no theoretical evidence exists so far. On the other hand, a recent, Monte-Carlo evaluation of $<\psi\bar{\psi}>$ in Schwinger's proper-time formalism, suggests that chiral symmetry is broken in the presence of intense EM fields which resemble those present at the GSI experiments [6]. Even though, there is no understanding so far of what causes this chiral transition, nor is it known, which is the effective coupling that may become critical in the heavy-ion system, there is some evidence supporting the view, that the relevant parameter here is a function of the background-field intensity [6,11], which depends on the charge, the energy and the distance of closest approach of the ions.

If this hypothesis is correct there are two points to be made. First, the intensity of the background EM-field is largest in so-called coherent collisions, in which the nuclei do not break up, i.e. for large impact parameters $b \geq R$, R being the nuclear radius, or equivalently for small momentum transfers $q \leq 1/R$. For heavy ions this corresponds to a maximum momentum transfer of 25-30 MeV, which is well above the $e+e-$ mass threshold, but very low compared to the energy of the ions , even at the GSI experiments. Therefore, if the anomalous peaks are genuine signatures of a phase transition, or some other collective excitations of the EM field, one would expect this to be an infrared, not an ultraviolet, effect, possibly due to the strong coupling of $e+e-$ pairs to the photon field induced in coherent heavy-ion collisions. The second point is that, once the idea of the formation of a supercritical vacuum is abandoned, there is no apriori reason why the mechanism which leads to the narrow peaks at GSI energies should not be effective also at high energies, were one may gain a better understanding of what is going on. In the following we shall explore this possibility, and its consequences for future heavy-ion experiments, in more detail.

The study of coherent heavy-ion collisions at relativistic energies has been mainly motivated by the possibility of discovering the Higgs boson and other heavy particles in such processes at future hadron colliders [12]. It was then noticed [13] that the cross section for producing a light system, like an $e+e-$-pair, via two photons emitted by the nuclei, is huge, of the order of a few kb's at 10 GeV/n *c.m.* energy, violating unitarity, which indicates the breakdown of the perturbation series in terms of α. The reason for this seems to be the "infrared" character of the e^+e^--pair emission which at high energies becomes analogous to the soft photon radiation.

The comparision of the two processes becomes particularly simple, if one adopts the semiclassical description for high-energy electromagnetic processes, whereby the EM field of a particle with charge Z and energy $E = \gamma M$ (in some reference frame) resembles a beam of real photons. The number of photons at an impact b with an energy ω is then given by [14]:

$$f(\omega, b) = \frac{Z^2\alpha}{\pi^2\beta^2b^2\omega}\,(\frac{\omega b}{\gamma\beta})^2\,[K_1^2(\frac{\omega b}{\gamma\beta}) + \frac{1}{\gamma^2}K_0^2(\frac{\omega b}{\gamma\beta})]\,, \tag{2}$$

and reduces to the much simpler expression: $Z^2\alpha/\pi^2b^2\omega$ for $\gamma = (1-\beta^2)^{-1/2} \gg 1$ and $\omega b/\gamma \leq 1$. Integrating this over the transverse impact plane d^2b with a lower cut-off $b_{min} = R$ yields the logarithmically divergent probability for soft photon emission from the nucleus:

$$\rho_\gamma \simeq \frac{2Z^2\alpha}{\pi}\int_{\omega_{min}}^{\gamma/R}\frac{d\omega}{\omega}\,ln(\frac{\gamma}{\omega R})\,. \tag{3}$$

A lower limit for the photon energy is usually set by the experimental energy resolution δE, while the infinite parts from real- and virtual-photon contributions with an energy below δE cancel, as it is well known, to each order of the coupling constant [15]. Nevertheless, the corresponding cross section for the emission of a single photon increases as the energy resolution increases and at some point violates unitarity. This is due to the fact that an accurate measurement of the photon energy implies a big uncertainty in the number of emitted photons, while perturbation theory describes processes which either increase or decrease the number of photons by one. Therefore, when the average number of photons ρ_γ is large and uncertain, the use of the coherent-state basis is more adequate, and eq.(3) is replaced by the probability of detecting any number of photons $\propto 1 - e^{-\rho_\gamma}$ with energy $\geq \delta E$ [16]. The exponential, which is obtained by summing over all photons that remain undetected, gives the probability for the vacuum-to-vacuum transition.

Now, in the equivalent photon approximation (EPA) [14], the differential cross section for a two-photon process factorizes into a part which contains the probability of emitting two almost real photons L, and the cross section for the underlying two-photon fusion process $\sigma(\gamma\gamma \to X)$:

$$d\sigma = \int\int L(\omega_1, \omega_2)\,\sigma_{\gamma\gamma\to X}(W^2)\,\delta(W^2 - 4\omega_1\omega_2)\,d\omega_1\,d\omega_2\,. \tag{4}$$

The two-photon luminosity function L, is obtained by integrating over the equivalent photon spectra of ion one and two over the impact $(\vec{b_1}\vec{b_2})$ plane, and the photon energies [17]:

$$L = \int\int_{b_1,b_2 \geq R_1,R_2} f_1(\omega_1, b_1)\, f_2(\omega_2, b_2)\, \Theta(B - R_1 - R_2)\, d^2b_1\, d^2b_2\,, \qquad (5)$$

where $B = |\vec{b_1} - \vec{b_2}|$ is the ion separation. Nuclear coherency is normally guarranteed by the step function Θ, but since the validity range of the EPA extends only where $q^2/m^2_{e^+e^-} \leq 1$, in this case only distant collisions with $B \geq 386$ fm can be treated reliably. Even at such large impact parameters the probability for producing an e^+e^- pair coherently [13],

$$\rho_{e^+e^-} = \frac{d^2\sigma_{e^+e^-}}{d^2B} \simeq \frac{Z_1^2 Z_2^2 \alpha^4}{\pi^2 m_e^2 B^2} \ln^2 \frac{2\gamma^2}{B m_e}\,, \qquad (6)$$

exceeds unity, *e.g.* in a lead-on-lead collision, already at c.m. energies of 100 GeV/n, which will be achieved in the near future at RHIC, the relativistic heavy-ion collider at Brookhaven.

The right way to think of this problem may be, that as the electron mass becomes negligible compared to the nuclear energy, the emission of e^+e^- pairs follows similar patterns as in the case of photons. Then, using a quasi-boson approximation, which is justified as long as the number of pairs emitted within a Compton volume is small, the expression in eq.(6) can be interpreted as the average number of pairs produced at a fixed impact parameter B, which at high energies becomes larger than one. So, it seems, that at high energies, multiple e^+e^- production is more favoured than the production of a single pair in a coherent heavy-ion collision. This interpretation is also in agreement with the leading behaviour of the elastic scattering amplitude of perturbative QED at asymptotic energies, given by Cheng and Wu [18],

$$T(s,t) \sim i\frac{2s}{M^2}\int d\vec{B}\, e^{i\vec{q}\vec{B}}\,(1 - e^{-\mathcal{A}})\,, \qquad (7)$$

where s and t are the c.m. energy and the momentum transfer squared. The amplitude $\mathcal{A}$ is the sum over diagrams which contain a tower of two photons connecting a number of e^+e^--loops to the field of the nuclei (see Fig.1). Though $\mathcal{A}$ grows as $s\ln^n s$, in the sum over all n-tower amplitudes $\mathcal{A}^n$ ($n = 1, ...\infty$), shown in Fig.1, unitarity is restored.

Since in heavy-ion scattering each tower is dominated by just one e^+e^- loop, so that $\mathcal{A}$ effectively reduces to $\rho_{e^+e^-}$ in eq.(6), by the optical theorem, the probability of emitting any number of pairs at a fixed impact parameter is $\simeq 1 - e^{-\rho}$, in agreement with our suggestion. Notice also that this quantity is the opacity of an almost black disc characterizing the absorption of the scattering amplitude which takes place at a certain energy and impact, due to the creation of low-energy e^+e^- pairs that get lost in the beam. If for a fixed energy, B is large, so that $\rho_{e^+e^-}$ is small, then $1 - e^{-\rho} \simeq \rho_{e^+e^-}$, which brings us back to the result of lowest order perturbation theory. On the other hand in the limit of complete absorption, *i.e.*, when the exponential goes to zero, the disc becomes opaque, saturating the Froissart bound for a strongly absorptive potential with a coupling constant which increases with the energy.

Up to now we have considered relativistic energies $\gamma > 1$ and large impact parameters $B \geq 1/m_e$ only. In order to make contact with the GSI data one has to extrapolate down to $\gamma \simeq 1$ and consider also impact parameters $R_1 + R_2 \leq B \leq 1/m_e$. By doing so one realizes two things. First of all, the non-perturbative character of the e^+e^- emission persists also for non-relativistic energies, and secondly, for impacts $B \leq 1/m_e$, eq.(7) cannot provide anymore the right answer to the problem of multiple e^+e^- production. To see this one may rewrite eq.(6) in terms of the effective coupling $g_{eff} \simeq Z_1 Z_2 \alpha^2 ln(2\gamma^2/x)$ of the e^+e^- pair to the background EM field, corresponding to the diagram of Fig.2a, and the scale $x = m_e B$, which measures the impact parameter in terms of the electron's Compton wavelength:

$$\rho_{e^+e^-} = \frac{g_{eff}^2}{\pi^2 x^2}. \tag{8}$$

Then, when $g_{eff} \geq 1$, so that multiple pair creation takes place, and $x \leq 1$, in which case Pauli's exclusion principle comes fully into play, the assumption, that the different e^+e^- pairs are created independently of each other, which lead to the exponentiation in eq.(7), is not sustainable. Instead, one would expect some new phenomenon to set in when this happens.

The velocity of the ions at the GSI experiments is $\beta \simeq 0.07$ and, for the typical range of measured angles, the minimal impact is $B_{min} \simeq 20$ fm. If one could use the simple expression for g_{eff}, which has been derived in the EPA, also for the nonrelativistic case, one would find that for all the nuclei that have been used at GSI (from $^{79}Au_{197}$ to $^{92}U_{238}$) the effective coupling is of order one. However, the validity of the EPA has been checked only for radiation processes in relativistic collisions, though other authors [19] have used it also for computing the two-photon production of e^+e^- pairs at GSI energies. Nevertheless, at least qualitatively, the functional dependence of g_{eff} on the nuclear charges, energy, and minimal impact should be the same also at low energies, and the evaluation of vacuuum polarization diagrams, coupled to the background field as shown in Fig.2b, seem to indicate that it is indeed of order one [6]. With this in mind we will now attempt a qualitative interpretation of the anomalies in the e^+e^- spectrum of heavy ions from low to high energies and discuss the consequences for future experiments.

There are two conceivable "scenarios" of how a system with n-pairs created within a Compton volume could respond. The first consists in a delocalization of the wave function of the electrons (and positrons) in ordinary space, which is equivalent to a correlation in momentum space. This could mean that the system originally formed within a space region with radius $r = B/2 \simeq (2m_e)^{-1}$ expands to a region with radius r^*, from where the e^+e^- pairs are finally emitted. As a matter of fact the GSI $e^\pm$ spectrum is characterized by two scales, which differ by an order of magnitude, or more. One of them is the time for which there is a substantial acceleration of the ions moving along the Rutherford trajectories with relative velocity β_0 [11]:

$$t_R \simeq \frac{10 Z_1 Z_2 \alpha}{\beta_0^3}\left(\frac{1}{M_1} + \frac{1}{M_2}\right) \simeq \frac{3}{m_e} \simeq 2 \cdot 10^{-21} s\,, \tag{9}$$

and sets the scale for the induced (*i.e.* the usual) e^+e^- pair production, which agrees with the measured width of the broad $e^\pm$ spectrum, $\Gamma_{e^+e^-} \simeq 500$ keV. The other scale is set by the intrinsic width of the narrow peaks, $\Gamma_{e^\pm} \leq 20 - 50$ keV, which lie superimposed on the broad spectrum of the induced production, and corresponds to a much longer lifetime:

$$t_{e^\pm} \leq 10^{-20} s \simeq \frac{10}{m_e}\,. \tag{10}$$

So the spacetime region from where the correlated electrons and positrons are emitted is by a few Compton wavelengths larger than the region that has been originally formed.

The other, somewhat "exotic" scenario, could involve a change of the intrinsic properties of the electron itself, like *e.g.* an increase of its mass to m_e^*, or equivalently a reduction of the Compton wavelength to λ_e^* and could therefore lead to the creation of tightly-bound, and therefore more localized, states of electrons, positrons and photons within the original formation volume. This is in fact what one would expect in the case of chiral symmetry breaking and the transition to a new phase of QED. Some interpretations of the GSI data have indeed used the analogy of positronium as an e^+e^- bound state of normal QED to fit the observed peaks to a spectrum of bound states of this postulated new phase [5]. Notice that according to our interpretation of eq.(8), both scenarios involve a change of the scale x, from $x \leq 1$ to $x \geq 1$, when the critical value of $g_{eff} \simeq 1$ is reached and $B \leq 1/m_e$.

At the moment, both scenarios are not yet theoretically founded, and are therefore a matter of speculation. However one may argue that the contribution of diagrams, where a large number of e^+e^- real and virtual pairs become coupled to each other via the background EM field of the ions, could lead to correlation or condensation phenomena. For example, the momenta of photons belonging to different towers in Fig.1 may be not independent. In fact, these are the only type of diagrams, that have been neglected in deriving eq.(7), which might be relevant in this case. Other diagrams, like the ones that involve the coupling of 2n external photons to each e^+e^- pair , as in Fig.2b, are less dominant, because, after all, each photon contributes an additional $Z\alpha$ factor to g_{eff}, which even for heavy nuclei is less than 0.7. Notice that g_{eff} can become of order one only because the logarithm from the fermion-loop integration can compensate the small deviation of $(Z\alpha)^2$ from unity. On the other hand, the summation of diagrams, which involve the exchange of an arbitrary number of photons between the nuclei, leads, in the eikonal approximation [16], also to exponentiation, and contribute to eq.(7) in a multiplicative way. Finally, diagrams including radiative corrections that do not exponentiate give constant contributions to other processes as well [15], and cannot be held responsible for the particularities of the e^+e^- emission.

Faced with various alternatives of how the presence of a nonperturbative effective coupling will affect the e^+e^- spectrum in heavy-ion collisions, it is important to think of ways to dinstinguish between them experimentally. In particular, because in relativistic collisions there is enough energy available to produce a large number of e^+e^- pairs in a single collision, one might be able, by measuring the correlations (or anticorrelations) between the momenta of different electrons and positrons, to draw some conclusions about the nature of the underlying production mechanism. This method, first known for photons as HBT (Hanbury-Brown-Twiss) interferometry, has become in connection with pions one of the main tools for probing the size of the space-time volume where a quark-gluon plasma, or a hadron gas, is thought to form when two relativistic nuclei collide at almost zero impact parameter [20].

The two-body correlation function $\mathcal{R}$, due to the interference of identical particles at small relative momenta q, is namely related to the size of the region in which they have been generated [21]:

$$\mathcal{R} = N \frac{d^6\sigma/d^3k_1 d^3k_2}{(d^3\sigma/d^3k_1)(d^3\sigma/d^3k_2)} \propto 1 \pm F^2(q^2). \tag{11}$$

In eq.(11), k_1 and k_2 are the 4-momenta of the two particles (*e.g.* $\pi^-\pi^-$, or, as in our case, e^-e^- and e^+e^+), so that, $q = k_1 - k_2$, and $F(q^2)$ is the Fourrier transform of the density distribution of the sources. N is determined by normalizing the single- and double-inclusive cross sections to the first and second binomial moments of the multiplicity distribution. The plus sign corresponds to the Bose-Einstein, the minus sign to the Fermi-Dirac interference. In deriving the right-hand side of eq.(11) a couple of assumptions have been made. First, it is assumed that the sources are statistically distributed and their internal structure is smooth and small in size compared to the total formation volume. Second, other correlations, *e.g.* due to final state interactions, have been neglected.

Let us now assume that the e^+e^- pairs, that will be measured in a heavy-ion collision for which nuclear fragmentation has been excluded and therefore the condition $B \geq R_1 + R_2$ is satisfied, are the decay product of some tightly-bound states formed, according to the second scenario, during the transition to a new (chiral) phase of QED. Because the time duration of the EM pulse, which represents the field of the moving ion, scales as $\Delta t \simeq B/\gamma \simeq 10^{-21}/\gamma\ s$, in relativistic collisions, the longitudinal dimension of the formation volume is small and so these new states will be produced almost instantaneously within a disc $\simeq \pi r^2$ with $(R_1 + R_2) \leq r \leq (2m_e)^{-1}$, provided that $g_{eff} \simeq 1$. The time component of the space-time distribution of the sources will then be dominated by the lifetime of the sources, $\tau \sim t_{e^\pm}$. Assuming for simplicity that the sources fill uniformely the disc and that they are all created simultaneously and decay according to an exponential law with mean time τ, we can parametrize the formfactor and therefore the correlation function of eq.(11) as follows:

$$\mathcal{R} \propto 1 + \frac{e^{-q^2 r^2/4}}{1 + q_0^2 \tau^2}, \tag{12}$$

where q_0 is the energy of the electron or positron. Notice that, because the e^+e^- pairs in this scenario are supposed to be the decay product of tightly-bound states, they will behave almost like bosons.

On the other hand, if the electrons and positrons are emitted, according to the first scenario, from a collective excitation of the vacuum polarization tensor, Bose-Einstein correlations will be absent, though, of course other type of momentum correlations might be present. [1] This is in general true for the emission of any number of particles belonging to the same quantum state and in particular, for bosons belonging to a coherent state. Therefore, e^+e^- pairs emitted within the conventional perturbative approach, eq.(7), will not exhibit any momentum correlations. This is mainly the reason why one may hope that the measurement of a correlation between electrons, and/or positrons, could one day be considered as a signal for a QED phase transition. There is of course a limit to the size that one can determine in this way, and for systems where $r \sim (2m_e)^{-1}$ this might be an impossible task, so one will have to choose a system for which g_{eff} becomes of order one only at a much smaller impact parameter, $B \simeq R_1 + R_2$. On the other hand, by triggering only on events where nuclear fragmentation is excluded, one can in principle avoid

[1]This possibility is presently under study.

the background from hadronic sources emitting e^+e^- pairs. For this a more detailed analysis will be needed.

Finally one should point out that with increasing multiplicity there will be a strong monochromatization and angular collimation of the pairs that form the interference maximum, and narrow "jets" of electrons and positrons with approximately equal energies will be produced [21]. As far as the modification of eqs.(11-12) due to final-state interactions is concerned, it is interesting to notice that, since the pairs are emitted at a time when the ions are so far apart that the EM field within the Compton volume is vanishing, eqs.(9-10), there will be no rescattering in the background field. So the only contribution comes from the interaction between the electrons and positrons themselves and can be taken into account just as in the case of pion emission.

1 Acknowledgements

I would like to thank L. Stodolsky for stimulating my interest in this problem and his contributions to the discussion.

Figure Captions

Figure 1. The contribution of a multi-tower diagram to the imaginary part of the elastic scattering amplitude.

Figure 2. The diagrams that give the effective coupling of an e^+e^- pair to the background EM field in coherent heavy-ion collisions. (a) The dominant contribution. (b) Higher order contributions.

References

[1] J. Schweppe *et. al.*, Phys. Rev. Lett. **51** (1983) 2261;M. Clemente *et. al.*, Phys. Lett. **B137** (1984) 41;T. Cowan *et. al.*, Phys. Rev. Lett. **54** (1985) 1761;*ibid.*, **56** (1986) 444; H. Tsertos *et. al.*, Zeit. Phys. **A326** (1987) 235;W. König *et. al.*, Zeit. Phys. **A328** (1987) 129;Phys. Lett. **B218** (1989) 12;M. Salabura *et. al.*, Phys. Lett. **B245** (1990) 153.

[2] Ya. B. Zeldovich and V. S. Popov, Usp. Fiz. Nauk **105** (1971) 403(Sov. Phys. Usp. **14** (1972) 673); B. Müller, J. Rafelski and W. Greiner, Zeit. Phys. **257** (1972) 62;*ibid.*, **257** (1972) 183; *see also*, W. Greiner, B. Müller and J. Rafelski,

Quantum Electrodynamics of Strong Fields (Springer Verlag, Berlin, 1985), and references therein.

[3] J. Rafelski, B. Müller and W. Greiner, Zeit. Phys. **A285** (1978) 49;Phys. Rev. **A24** (1981) 103;J. Rafelski, U. Müller, B. Müller and W. Greiner, Zeit. Phys. **A303** (1981) 173;U. Müller *et. al*, Zeit. Phys. **A313** (1983) 263.

[4] For a review, see J. van Klinken, in Proceedings of the XXIVth Rencontres de Moriond, Les Arcs, Jan. 1989.

[5] For a review, see D. G. Caldi, Comments Nucl. Part. Phys. **19** (1989) 137,and references therein.

[6] D. G. Caldi, in Proceedings of the NATO Advanced Study Institute on Vacuum Structure in Intense Fields, Cargese, France, July 1990; D.G. Caldi and S. Vafaeisefat, *Background electromagnetic fields induce phase transition to new QED*, preprint in preparation.

[7] L. D. Landau, in Niels Bohr and the Development of Physics, McGraw Hill, New York, 1955.

[8] E. Dagotto and J. Kogut, Phys. Rev. Lett. **59** (1987) 617;Nucl. Phys. **B295** (1988) 123;J. Kogut, E. Dagotto and A. Kocic, Phys. Rev. Lett. **60** (1988) 772.

[9] T. Maskawa and H. Nakajima, Prog. Theor. Phys. **52** (1974) 1326;*ibid* **54** (1975) 860; R. Fukuda and T. Kugo, Nucl. Phys. **B117** (1976) 250;V.A. Miransky, Il Nuovo Cim. **90A** (1985) 149; V.A. Miransky and P.I. Fomin, Sov. J. Part. Nucl. **16** (1985) 203.

[10] C.N. Leung, S.T. Love and W.A. Bardeen, Nucl. Phys. **B273** (19868) 649;V.A. Miransky, Phys. Lett. **B248** (1990) 151.

[11] R.D. Peccei, J. Sola and C. Wetterich, Phys. Rev. **D37** (1988) 2492.

[12] E. Papageorgiu, Phys. Rev. **D40** (1989) 92;Nucl.Phys. **A498** (1989) 593c;M. Grabiak, B. Müller, W. Greiner, G. Soff and P. Koch, J. Phys. **G15** (1989) L25;M. Drees, J. Ellis and D. Zeppenfeld, Phys. Lett. **B223** (1989) 454;K.J. Abraham *et.al*, Proc. of the Large Hadron Collider Workshop, ed. G. Jarlskog and D. Rein, Aachen, 1990.

[13] E. Papageorgiu, Ph.D. Thesis and MPI Report **68** (1989) ;G. Baur, talk given at the CBPF Int. Workshop on Relativistic Aspects of Nuclear Physics, Rio de Janeiro, 1989.

[14] See *e.g.*, J.D. Jackson, Classical electrodynamics (Wiley, New York, 1963); V.M. Budnev, I.F. Ginzburg, G.V. Meledin and V.G. Serbo, Phys. Rep. **C15** (1975) 181.

[15] D.R. Yennie, S.C. Frautschi and H. Saura, Ann. Phys. **13** (1961) 379.

[16] See *e.g.*, C. Itzykson and J.B. Zuber, Quantum Field Theory (McGraw-Hill, 1980).

[17] E. Papageorgiu, Phys. Lett. **B250** (1990) 155.

[18] H. Cheng and T.T. Wu, Phys. Rev. Lett. **24** (1970) 1456.

[19] G. Alexander, E. Gotsman and U. Maor, Zeit. Phys. **C34** (1987) 329.

[20] For a review see, H. Satz in Proc. of the Large Hadron Collider Workshop, ed. G. Jarlskog and D. Rein, Aachen, 1990.

[21] M. Gyulassy, S.K. Kauffmann and L.W. Wilson, Phys. Rev. **C20** (1979) 2267;M.I. Podgoretskii, Sov. J. Part. Nucl. **20** (1989) 266,and references therin.

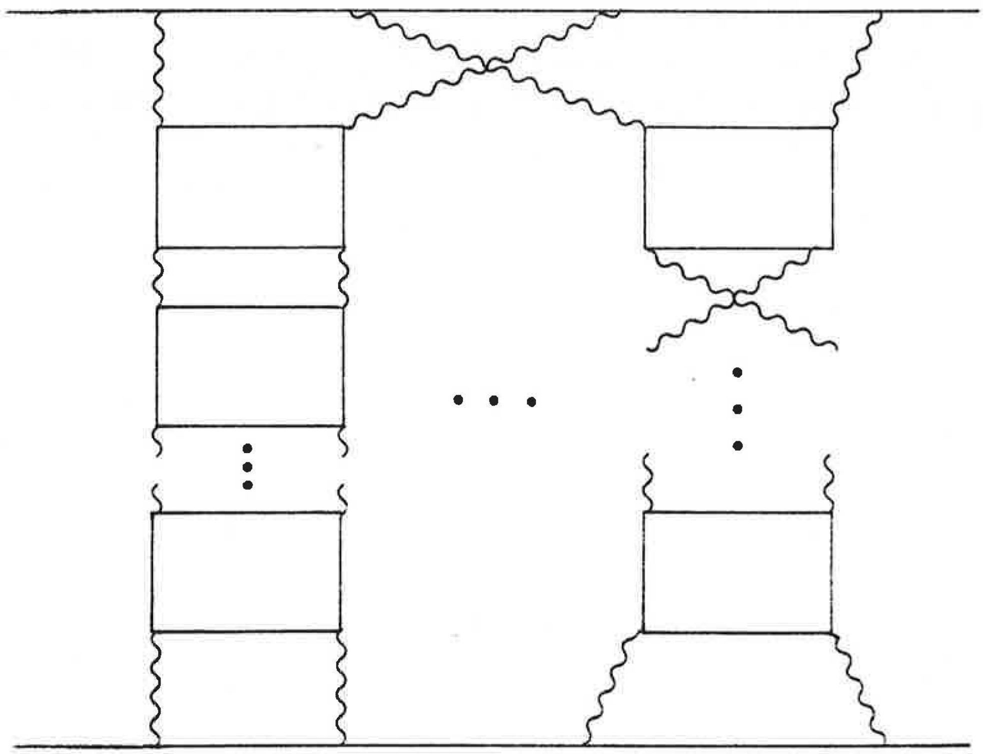
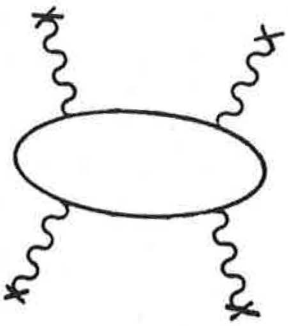
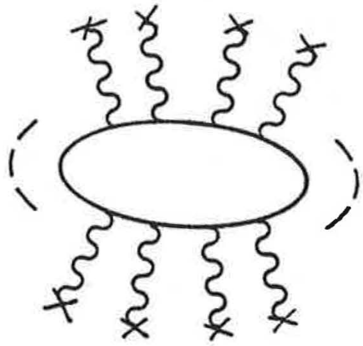


Fig. 1



(a)



(b)

Fig. 2