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ABSTRACT

The status of order parameters in lattice gauge theories with matter fields is reviewed, together with several related problems. Then the flux correlations order parameter is discussed in detail. An analysis of its finite distance behaviour leads to the definition of a characteristic length scale r_D. This is possible both at zero and at finite temperature. It is argued that r_D is the screening length for dynamical charge fluctuations (Debye screening). In general, the length defined from the exponential decay of Polyakov loop correlations is different from r_D.

1. INTRODUCTION

In this talk I will try to present the current status of the order parameters (OP) proposed some time ago by Klaus Fredenhagen and myself.

Typically a lattice gauge theory (LGT) with matter fields has three regions: a free charge phase, a screening (Higgs) region and a confinement region /1, 2/. In an actual model some of these regions may be absent. If the matter fields are in the fundamental representation of the gauge group, the Higgs and the confinement regions are not separate phases.

For a long time after LGT's were invented, no OP was known except in the limiting cases of pure matter and pure gauge theories /3/. In ref. /4/ we proposed three different OP's. In fact the vacuum overlap OP (VOOP) had been proposed as early as 1982 /5/.

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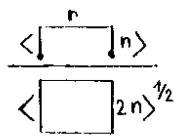
integral formulation of the model is:

$$\rho(r,n) = \langle 0 | \underline{x}, \underline{y}, n \rangle = \tag{2.2}$$



Here $r = |\underline{x}-\underline{y}|$, and the symbolic notation means a gauge invariant path ordered product of the gauge fields along the open string and the closed loop resepctively.

The VOOP is the $r,n \to \infty$, $n \gtrsim r$ limit of $\rho(r,n)$. In the free charge phase the limit should be zero, since a charged state is orthogonal to the vacuum. In the Higgs and confinement regions the limit should be nonzero. Thus the criterion for existence of charged states reads:

$$\lim_{\substack{r,n\to\infty \\ r \lesssim n}} \rho(r,n) = \begin{cases} = 0 & \text{free charge} \\ \neq 0 & \text{confinement, Higgs} \end{cases} \tag{2.3}$$

For Abelian gauge groups and scalar matter fields, eq. (2.3) follows immediately from Griffith inequalities. In models with Fermions or a non-Abelian gauge group, (2.3) has not been proven, with the exception of the massless Schwinger model (in the continuum!) /11/.

The construction of charged states in Z_2 theories, which was our starting point in /4/, can be easily generalized only for models having a convergent expansion in the free charge region. Unfortunately this excludes most cases of real interest.

In the Z_2 model even more can be rigorously established. Barata and Fredenhagen proved that the charged states are really particle states /12/.

A totally different construction of the charged states in Z_2 theories was given by Szlachányi /13, 14/. He formulates in a gauge invariant way the old idea /2/ that the charged fields can be constructed using the matter fields in an appropriately chosen gauge. Then he succeeds in finding a suitable class of gauges and goes on to construct the charged states. The resulting charged sector is identical to that obtained by our construction. Moreover, his proofs

The VOOP has the most rigorous foundation of our OP's. It is based on the following ideas for constructing a charged state /4, 6, 7/: separate a charge-anticharge pair, regularize the energy, then send the anticharge to infinity, and check whether the resulting state is indeed charged. The quantity testing whether this candidate for a charged state is indeed charged is its scalar product with the vacuum (the "vacuum overlap"). In the limit of the pure matter theory the VOOP goes into the square of the usual local order parameter (e.g. magnetization). In the Higgs region the VOOP is the "Higgs expectation value", defined properly, in a gauge invariant way.

The construction of charged states /4/ that led to the VOOP was also meant as an example for the general analysis of particle states in theories with a local gauge symmetry /8/. This general analysis also suggested to investigate correlations of electric fluxes associated with the gauge group center, in order to understand those properties of the vacuum which are responsible for the existence or absence of charged states. For theories with a mass gap, a useful quantity for studying these correlations is the flux correlations OP (FCOP), defined in ref. /4/. In the limit of the pure gauge theory, the FCOP goes into the 't Hooft loop /9/ (the dual of Wilson loop in the case of selfdual theories). As opposed to the VOOP, the FCOP is well-defined even at finite temperature.

Our third OP compares the charge of the candidate for a charged state to that of the vacuum. We therefore call it the charge measurement OP (CMOP). In ref. /4/ this comparison was possible although we did not succeed in regularizing the cahrge operator itself. In this talk the CMOP will not be discussed further.

A discussion of the three OP's in terms of general properties of particle states in gauge theories can be found in ref. /10/. Here the accent will lie on the ideas underlying the OP's and their practical use in LGT's.

All our OP's are limits of local quantities. Before the limit, these quantities contain valuable additional information. In pure gauge theories, knowledge of the finite Wilson loops allows us to determine the quark-antiquark potential. In ref. /7/ it was shown that in the confinement region of LGT's with matter fields, the VOOP before the limit can be used to define a length scale. It was argued that this is the scale where quark fragmentation sets in. In this talk it will be shown that the FCOP also allows us to define a length scale, this time in the Higgs region. Furthermore, at finite temperature this length scale may be defined for all values of the coupling constants. It will be argued that it can be interpreted as the screening length for dynamical charge fluctuations (Debye screening).

All other OP's for LGT's with matter fields that have been proposed in the literature and that correctly reproduce the known phase diagrams, are closely related to the VOOP. This will be discussed in section 2, in a brief survey on the work done on the VOOP and related problems. In section 3 the FCOP will be presented and the model of a plasma for the vacuum or for a thermal equilibrium state will be discussed.

For simplicity, the discussion will be given an terms of a LGT with a scalar matter field in the fundamental representation of the gauge group. The action is

$$S = -\beta \sum_p \chi(U(p)) - \kappa \sum_l 2\,\mathrm{Re}(\varphi^\dagger U \varphi)(l) + \sum_x V(\varphi(x)) \qquad (1.1)$$

Here p, l, x are the plaquettes, links and sites of a d-dimensional Euclidean lattice, U are the gauge fields, φ are the matter fields, and χ is some character that contains the fundamental character as an irreducible component. The time-zero field operators will be denoted by $\hat{U}(\underline{l})$ and $\hat{\varphi}(\underline{x})$. In general, underlined quantities will denote geometrical objects in d-1 (space) dimensions.

2. THE VACUUM OVERLAP ORDER PARAMETER AND RELATED PROBLEMS

A candidate for a charges state is obtained by considering the following sequence of norm-one vectors /4, 6, 7/:

$$|\underline{x},\underline{y},n\rangle = \frac{1}{\|\cdot\|} \sum_{a,b} \hat{\varphi}_a(\underline{x})\, \hat{\varphi}_b^+(\underline{y})\, \hat{T}^n\, \hat{U}(\underline{L})\, |0\rangle \qquad (2.1)$$

where a,b are gauge group indices, $\underline{L}$ is a spatial path connecting $\underline{x}$ and $\underline{y}$, $\hat{U}(\underline{L})$ is the path ordered product of $\hat{U}$'s along $\underline{L}$ and $\hat{T}$ is the transfer matrix (derived in the usual way in the temporal gauge). For simplicity, let $\underline{L}$ be along a coordinate axis. The Euclidean time translation by n achieves the regularization of the energy for the charge-anticharge pair: the relative weight of the low-energy components of $|\underline{x},\underline{y},n\rangle$ increases with n. Actually for an Abelian theory we can extend the methods used in /4/ to prove that the energy of $|\underline{x},\underline{y},n\rangle$ stays bounded in the limit $|\underline{x}-\underline{y}| \to \infty$, $n \to \infty$ provided n grows at least linearly with $|\underline{x}-\underline{y}|$ (this statement is true for the whole phase diagram). Unfortunately the proof cannot be easily generalized for non-Abelian theories.

Let us denote by $\rho(r,n)$ the vacuum projection of $|\underline{x},\underline{y},n\rangle$. Up to a minor modification /6, 7/, the expression of $\rho(r,n)$ in terms of expectation values in the d-dimensional Euclidean path-also rely heavily on the convergent expansion.

For QED we expect our charged state to be similar to the physical electron in the Coulomb gauge, i. e. the bare electron dressed with the electric field of a pointlike external source. The same is true for Szlachányi's construction, if he does it in a rotationally invariant way. In ref. /14/ he gives a method to generalize to arbitrary compact gauge groups the gauge invariant definition of expectation values of products of time-zero Coulomb-gauge matter fields.

A similar construction using the Landau gauge was proposed by Kennedy and King /15/. Borgs and Nill /16/ showed that, due to problems with Gribov copies, this construction in fact works only for the noncompact scalar lattice QED.

In the Higgs phase the Szlachányi and Kennedy-King constructions define the "Higgs expectation value" in a gauge invariant way. The same is true for the VOOP. Numerically, these different definitions will not be identical. However, this is not necessary for a correct discussion of the Higgs mechanism.

As opposed to the VOOP, the ideas of Szlachányi and Kennedy-King do not seem to work in the confinement region.

The behaviour of $\rho(r,n)$ for finite r and n was discussed in detail in ref. /7/. In the free charge phase and in the Higgs region, $\rho(r,\infty)$ decays exponentially with r to the asymptotic value of the VOOP. From this decay we can determine the charged particle mass and the Higgs mass respectively. This is however difficult numerically, since it is difficult to compute Euclidean correlations of zero momentum charged states. In the confinement region, $\rho(r,\infty)$ decays exponentially with the perturbative charged particle mass for small values of r, while for large values it shoots up again to the asymptotic value of the VOOP. The scale of the crossover is the scale where fragmentation of the charge-anticharge pair sets on. The asymptotic value of the VOOP can be estimated more easily using the Bricmont-Fröhlich parameter /17/, which is equal to the VOOP in the Higgs-confinement phase. In the free charge phase, the Bricmont-Fröhlich parameter actually tests the existence of a hydrogen atom /7, 4/.

The VOOP has been investigated numerically in a variety of Monte Carlo simulations. The most accurate study was possible in the Z_2 model in 4 dimensions /7, 18/, where the VOOP was used to give strong numerical evidence for a line of second order phase transitions with mean field exponents separating the free charge phase from the Higgs region. The Aachen-Georgia group have studied the VOOP in the SU(2) Higgs model in 4 dimensions /19/. This case

is much more difficult numerically than the Z_2 case. Within their accuracy, the predictions of /7/ were verified. This is very important, since the theoretical arguments are much weaker for non-Abelian gauge groups. The VOOP has also been computed in the 4-dimensional U(1) Higgs model. For matter fields with unit charge, the Aachen group has investigated the Higgs-free charge transition /20/. They have some evidence that the transition is first order. The case of doubly charged matter fields was simulated by Azcoiti and Tarancón /21/. Alessandrini et al. /22/ analized it using mean field methods. The U(1) numerical studies also confirm the theoretical predictions for $\rho(r,n)$.

The VOOP cannot be defined at a finite temperature T since the lattice is finite in Euclidean time. At finite T the theory effectively decouples (see e.g. /23/) into a pure matter theory with magnetic field and a LGT with matter fields, both in d-1 dimensions and at zero temperature. If we consider the paths of eq. (2.2) as lying in a specelike plane, $\rho(r,n)$ will define the VOOP for the d-1 dimensional LGT with matter fields. It can be used to study e.g. the "symmetry restauration" transition, which was repeatedly discussed in this conference.

3. THE FLUX CORRELATION ORDER PARAMETER AND THE SCREENING OF DYNAMICAL CHARGE FLUCTUATIONS

The general analysis of charged states in a massive gauge theory with matter fields /8/ leads to the following picture. A charged state can be constructed from a charge-anticharge pair by sending the anticharge to infinity. Things can always be arranged such that asymptotically the lines of electric flux starting at the charge localization point and going to infinity are inside a cone of given solid angle. The total charge can be determined by measuring the electric flux through an arbitrarily large closed surface (Gauss' law). However, since the asymptotic direction of the cone axis is not an observable, it is not possible to measure a nontrivial electric flux through a solid angle of less than 4π. This latter property prevents the existence of charged states carrying an additive charge (Swieca's theorem) /24/. On the other hand, states carrying a multiplicative charge may exist if electric fluxes in different directions are correlated strongly enough.

This general picture is not contradicted by any LGT study in the literature. In a LGT Gauss' law holds only for the center $\mathcal{C}$ of the gauge group. Thus a charged state may only carry the quantum numbers of $\mathcal{C}$. In all known examples of a massive free charge phase (the photon is massive) $\mathcal{C}$ is a finite Abelian group, and therefore the charge is multiplicative rather than additive.

Before defining the FCOP some additional notation must be introduced. Let $\hat{\mathcal{E}}_V(\underline{l})$ be the left translation operator with the element V of the gauge group. As in the case of the gauge fields, it is useful to consider $\underline{l}$ as an orient ed link. If $\underline{l}$ and $\underline{l}'$ differ only by their orientation, then $\hat{\mathcal{E}}_V(\underline{l}') = \hat{\mathcal{E}}_{V^+}(\underline{l})$. For a set $\underline{\mathcal{L}}$ of oriented links and an element C of $\mathcal{C}$, let us define $\hat{\mathcal{E}}_C(\underline{\mathcal{L}}) = \prod_{\underline{l}\in\underline{\mathcal{L}}} \hat{\mathcal{E}}_C(\underline{l})$. Then

$$\langle 0|\,\hat{\mathcal{E}}_C(\underline{\mathcal{L}})|0\rangle = \Big\langle \prod_{\underline{l}\in\underline{\mathcal{L}}} \exp \beta\{\chi(CU(p_{\underline{l}}))-\chi(U(p_{\underline{l}}))\}\Big\rangle \qquad (3.1)$$

Here $p_{\underline{l}}$ is a time-like plaquette between the time-zero and time-one hyperplanes such that $\underline{l}$ is the part of $\partial p_{\underline{l}}$ contained in the time-zero hyperplane.

Let $\underline{\Lambda}$ be a cube in d-1 dimensions with side length r. Let $\partial^*\underline{\Lambda} = \underline{\mathcal{L}} \cup \underline{\mathcal{R}}$ be the decomposition of the coboundary of $\underline{\Lambda}$ into a left and right "hemisphere" (all geometrical considerations are in d-1 dimensions; on the dual lattice, $\underline{\mathcal{L}}$ and $\underline{\mathcal{R}}$ are d-2 dimensional open surfaces with the same boundary). Denote by $\underline{\mathcal{M}}$ the minimal set of spatial links such that $\partial^*\underline{\mathcal{L}} = \partial^*\underline{\mathcal{R}} = \partial^*\underline{\mathcal{M}}$ (minimal surface on the dual lattice).

Assume now that we are in a massive free charge phase and we try to bring a charge inside $\underline{\Lambda}$ such that the electric flux coming in from infinity is localized in a cone. Since the asymptotic direction of the cone axis is not observable, there are vacuum fluctuations that delocalize the electric flux. The probability for the vacuum to contain large closed lines of electric flux is relatively high. As a consequence the vacuum expectation values of the electric flux operators $\hat{\mathcal{E}}_C(\underline{\mathcal{L}})$ and $\hat{\mathcal{E}}_C(\underline{\mathcal{R}})$ $(C\in\mathcal{C})$ are strongly correlated (it only makes sense to consider fluxes associated with the gauge group center). If on the other hand we are in a massive phase without charged states, the asymptotic direction of the cone could be observed by measuring the electric flux through appropriate portions of $\partial^*\underline{\Lambda}$. There are no vacuum fluctuations delocalizing the electric flux. This also means that in the vacuum the fluxes through $\underline{\mathcal{L}}$ and $\underline{\mathcal{R}}$ are only weakly correlated.

Since the electric fluxes are multiplicative, the quantity to be considered for a study of the flux correlations is:

$$f_C(\underline{\Lambda}) = \frac{\langle 0|\,\hat{\mathcal{E}}_C(\underline{\mathcal{L}})|0\rangle\,\langle 0|\,\hat{\mathcal{E}}_C(\underline{\mathcal{R}})|0\rangle}{\langle 0|\,\hat{\mathcal{E}}_C(\partial^*\underline{\Lambda})|0\rangle} \qquad (3.2)$$

If our intuitive picture of the role played by vacuum fluctuations is correct, in a massive free charge phase the main contributions to (3.2) will come from the closed lines of electric flux passing through both $\underline{\mathcal{L}}$ and $\underline{\mathcal{R}}$ (remember that $\hat{\mathcal{E}}_C(\underline{\mathcal{L}})$ and $\hat{U}(\underline{L})$ do not commute if $\underline{L}$ passes through $\underline{\mathcal{L}}$). The other vacuum fluctuations play a similar role in the numerator and denominator and we expect them to cancel up to a contribution at the "perimeter" $\partial^*\underline{\mathcal{L}} = \partial^*\underline{\mathcal{R}} = \partial^*\underline{\mathcal{M}}$. The closed lines passing through $\underline{\mathcal{L}}$ and $\underline{\mathcal{R}}$ can be classified according to their intersection with the minimal surface $\underline{\mathcal{M}}$. This immediately suggests that (3.2) decays with the area of $\underline{\mathcal{M}}$. In a phase without free charges this area effect is absent and (3.2) will decay with the perimeter of $\underline{\mathcal{M}}$. To sum up,

$$f_C(\underline{\Lambda}) \underset{\underline{\Lambda}\to\infty}{\sim} \begin{cases} e^{-c_1|\underline{\mathcal{M}}|} & \text{free charge} \\ e^{-c_2|\partial^*\underline{\mathcal{M}}|} & \text{confinement, Higgs} \end{cases} \qquad (3.3)$$

We call $f_C(\underline{\Lambda})$ the FCOP.

By Gauss' law, the denominator of (3.2) is nothing else than the vacuum expectation value of the operator measuring the charge inside $\underline{\Lambda}$. In the limit of the pure gauge theory it becomes 1. In this limit the numerator of (3.2) converges towards the square of the 't Hooft loop /9/, which is known to have the behaviour (3.3).

For the Z_2 theory (3.3) was proven using convergent expansion techniques /4/.

In the d=3 Z_2 theory the FCOP is dual to the VOOP. In particular, in the Higgs region the FCOP behaves differently for small and for large r, since it is dual to the VOOP in the confinement region. This situation is not restricted to d=3. For arbitrary dimensions, the 0th order contribution to $f_{-1}(\underline{\Lambda})$ in the Higgs region expansion is

$$f_{-1}(\underline{\Lambda}) = \frac{\left(e^{-c_2|\partial^*\underline{\mathcal{M}}|-2\beta|\underline{\mathcal{L}}|} + e^{-2\beta|\underline{\mathcal{M}}|-\varkappa|\underline{\Lambda}|}\right)^2}{e^{-2\beta|\partial^*\underline{\Lambda}|} + e^{-2\varkappa|\underline{\Lambda}|}} \qquad (3.4)$$

Thus for large values of β and $\varkappa$,

$$f_{-1}(\underline{\Lambda}) \sim \begin{cases} e^{-2\beta|\underline{\mathcal{M}}|} & r \ll r_D \\ e^{-c_2|\partial^*\underline{\mathcal{M}}|} & r \gg r_D \end{cases} \qquad (3.5)$$

where $r_D = 2d\beta/\varkappa$.

The scale r_D can be interpreted in the framework of a plasma model for the Higgs phase. The vacuum is a condensate of charged particles. For distances smaller than r_D, the flux correlations are strong. If by some fluctuation a charged particle came inside Λ , a charge measurement could be performed by measuring the electric flux through $\partial^*\Lambda$, provided the linear size of Λ is smaller than r_D. Thus r_D is the scale at which the fluctuations of the dynamical charge are screened. Usually this is called the Debye screening length /25/.

In Plasma Electrodynamics the screening length r_{pot} for the potential (i.e. for the energy of an external charge-anticharge pair) is identical to r_D. This is shown by using Gauss' law to relate the charge density and the potential /25/. However, a similar argument is not possible if Gauss' law holds only in multiplicative form, as in our case. A 0th order computation of r_{pot} in the Higgs region of the Z_2 model results in:

$$r_{pot} = \frac{1}{4(d-1)\beta + \varkappa} \qquad (3.6)$$

In an electrodynamical plasma $r_D \to \infty$ as $\beta \to \infty$ (zero gauge coupling). The same is true in the Z_2 model for r_D but not for r_{pot}.

The FCOP can be considered also at finite temperatures T, since eq. (3.2) involves only time-zero operators. The Debye screening length r_D can now be defined in the whole phase diagram, not only in the Higgs region. We expect the following qualitative behaviour:

a) If at T = 0 we are in the free charge phase, at finite T we will have a plasma of charged particles. r_D diverges as $T \to 0$ and monotonically decreases as T grows. As $T \to \infty$, r_D settles to a value around one lattice spacing.

b) If at T = 0 we are in the Higgs region and $r_D > 1$, r_D will monotonically decrease with increasing T. The asymptotic value is again of the order of 1.

c) If at T = 0 we are in the confinement region, then we expect r_D to be very small as T increases towards the deconfinement transition T_d. At T_d there is no phase transition, but a crossover from the confinement regime to a charged particle plasma regime. r_D will rise sharply as T reaches T_d from below, and then decrease as described in a) - b) if T is increased further.

For the Z_2 model simple arguments can be given that support the finite T behaviour of r_D described here /26/. We also plan to

do a Monte Carlo investigation of the FCOP. In general the FCOP is difficult to simulate for large values of β, as seen from eq. (3.1). However, for the Z_2 case a duality transformation can be used such that in the new model one has to compute expectation values of products of $\pm$ 1.

The decay of Polyakov loop correlations can be used to compute r_{pot} at finite T. There is again no similarity at all between r_D and r_{pot}. However, it is quite plausible that the deconfinement peak of r_D coincides with the rapid rize in the value of one Polyakov loop. Unfortunately we have no good theoretical argument to support this belief.

Eq. (3.4) suggests that one can use the denominator of (3.2) alone in order to define r_D and to investigate the deconfinement transition.

The study of the FCOP is just at its beginning. Reliable methods to compute it in various models have to be developed. The theoretical foundations are also not as strong as for the VOOP. Maybe different quantities can be defined that have essentially the same meaning but are easier to calculate. It would be interesting to analyse the role played by flux correlations in massless phases. Last but not least, it would be interesting to see how far the plasma analogy can be pushed.

REFERENCES

1 K. Osterwalder and E. Seiler, Ann. Phys. 110 (1978) 440

2 E. Fradkin and S. Shenker, Phys. Rev. D19 (1979) 3682

3 K.G. Wilson, Phys. Rev. D10 (1974) 2445

4 K. Fredenhagen and M. Marcu, Commun. Math. Phys. 92 (1983) 81

5 K. Fredenhagen, Talk presented at the Colloquium in honour of Prof. R. Haag at the occasion of his 60th birthday, Hamburg, November 15, 1982, appeared as Freiburg University preprint THEP 82/9

6 K. Fredenhagen and M. Marcu, Phys. Rev. Lett. 56 (1986) 223

7 M. Marcu, in "Lattice gauge theory - a challenge in large-scale computing", B. Bunk, K.H. Mütter and K. Schilling editors, Plenum (1986) p. 267

8 D. Buchholz and K. Fredenhagen, Commun. Math. Phys. 84 (1982) 1

9 G. 't Hooft, Nucl. Phys. B153 (1979) 141

10 K. Fredenhagen, Lectures given at the Erice Summer School on "Fundamental Problems of Gauge Field Theory" (1985), DESY preprint 85-120

11 J.L. Alonso and A. Tarancón, Phys. Lett. 165B (1985) 167

12 J. Barata and K. Fredenhagen, "Charged particles in Z_2 gauge theories", to be published

13 K. Szlachányi, Phys. Lett. 147B (1984) 335

14 K. Szlachányi, Budapest Central Research Institute of Physics preprint KFKI-1986-35/A

15 T. Kennedy and C. King, Phys. Lett. 55 (1985) 776 and Commun. Math. Phys. 104 (1986) 327

16 C. Borgs and F. Nill, Nucl. Phys. B270 (FS16) (1986) 92

17 J. Bricmont and J. Fröhlich, Phys. Lett. 122B (1983) 73

18 T. Filk, K. Fredenhagen and M. Marcu, Phys. Lett. 169B (1986) 405

19 H.G. Evertz et al., Phys. Lett. 175B (1986) 335

20 J. Jersák, this conference,
Aachen group, in preparation

21 V. Azcoiti and A. Tarancón, Phys. Lett. 176B (1986) 153

22 V. Alessandrini et al., Zaragoza preprint DFTUZ-86.9

23 B. Svetitzky and L.G. Yaffe, Nucl. Phys. B210 (FS6) (1982) 423

24 J.A. Swieca, Phys. Rev. D13 (1976) 312,
D. Buchholz and K. Fredenhagen, Nucl. Phys. B154 (1979) 226

25 A.F. Alexandrov, L.S. Bogdankevich and A.A. Rukhadze, "Principles of Plasma Electrodynamics", Springer (1984)

26 K. Fredenhagen and M. Marcu, in preparation