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Polarization and Chirality: the quantum features of the Quark Gluon Plasma

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Abstract

Polarization and chirality are direct manifestations of quantum mechanics in the Quark Gluon Plasma as a relativistic fluid. This is one of the reasons why they are intriguing phenomena, that have attracted so much attention lately. In this talk I will review what is, in my view, the current theoretical understanding and highlight the most interesting recent results.

1. Introduction

Much effort has been recently devoted, in relativistic heavy ion physics, to the search of chirality and polarization-related effects. The Chiral Magnetic Effect (CME) is yet to be discovered in heavy ion collisions but the theoretical work is now entering a mature stage where the endeavour for a full numerical simulation including dynamical electromagnetic field is ongoing. On the other hand, polarization is, at least on the experimental side - certainly in a more advanced stage. The first positive evidence of a global polarization reported by the experiment STAR [1] has been confirmed by more and more accurate measurements that have been presented in this conference [2, 3]. This result has generated much enthusiasm not just among those working on the subject, but also among those interested in chiral- related effects. Indeed, these two phenomena have several common features: they are both related to the polarization degrees of freedom, they should arise only in peripheral collisions, and others. On top of that, they are direct manifestations of the quantum nature of the Quark Gluon Plasma (QGP); CME requires an anomaly at work, which is by definition a quantum breaking of a classical symmetry and polarization is itself a quantum object.

2. The Chiral Magnetic Effect

I will not bring up anything about the experimental search of the CME, particularly about the correlators designed to obtain an evidence thereof. For that topic and an experimental overview on chirality and polarization I refer the reader to the talk by Tu in this conference [4].

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The search for the CME is now entering - as has been mentioned in the Introduction - a more advanced stage where one aims at dynamically evolving both the axial charge and the electromagnetic field to make more accurate predictions. Supplementing the usual hydrodynamic model of QGP with these new quantities implies the transition from the currently available viscous hydrodynamics code to the so-called Chiral-Relativistic-Magneto-Hydrodynamics (χ-RMHD). To date, most studies have focused on the theoretical features of χ-RMHD [5, 6, 7] in terms of constitutive equations and their analytic expressions. In fact, no numerical implementation is available, although there are ongoing projects to develop it.

While the full χ-RMHD is yet to come, there have been worthwhile intermediate steps, like ideal RMHD codes [8, 9] as well as a hybrid code like AFVD (Anomalous Viscous Fluid Dynamics) [10] wherein vector and axial currents are dynamically evolved in a background electromagnetic field, while the fluid viscous hydrodynamics is completely decoupled. With these calculations, a charge imbalance for the final particles is obtained, which is of course strongly dependent on the initial value of the axial charge [11].

It is worth reviewing the physics of the CME from the perspective of equilibrium quantum statistical mechanics. Thermodynamic equilibrium in quantum statistical mechanics, in the grand-canonical ensemble, is defined through the density operator maximizing entropy $S = -\mathrm{tr}(\widehat{\rho}\log\widehat{\rho})$ with the constraints of given mean values of *conserved* quantities, such as energy, charge etc., which leads to a density operator of the form:

$$\widehat{\rho} = \frac{1}{Z}\exp[-\widehat{H}/T + \mu\widehat{Q}/T] \tag{1}$$

where $\widehat{H}$ is the Hamiltonian operator and $\widehat{Q}$ the conserved charge. The parameters $1/T$ and μ/T are - as it is known - the Lagrange multipliers associated to the constraints. For a system of free massless fermions, there is an additional conserved charge besides the electric/vector charge, namely the axial charge:

$$\widehat{Q}_A = \int \mathrm{d}\Sigma_\mu \, \widehat{j}^\mu_A \tag{2}$$

which is written here in a fully covariant form. The integration can be done over any space-like 3D hypersurface because the divergence of the integrand vanishes. Therefore, since the axial charge is conserved, one can write a density operator including an *axial chemical potential* μ_A:

$$\widehat{\rho} = \frac{1}{Z}\exp[-\widehat{H}/T + \mu\widehat{Q}/T + \mu_A\widehat{Q}_A/T] \tag{3}$$

A value $\mu_A \neq 0$ signals an imbalance between right-handed and left-handed fermions.

For the mean values of general physical quantities to be non-vanishing at equilibrium, the symmetries of the density operator (3) are crucial. In order for an electric current to be allowed at global thermodynamic equilibrium, rotational symmetry needs to be broken along with time reversal, charge conjugation and parity. For the former three, a constant and uniform magnetic field is enough as it breaks them (see table 1) in the Hamiltonian. Yet, a constant and uniform magnetic field does not break parity; for this purpose, a non-vanishing axial chemical potential - that is a chiral imbalance - is necessary. This is why in the renowned formula [12]:

$$\mathbf{j}_{el} = \frac{e^2}{2\pi^2}\mu_A\mathbf{B}$$

the coefficient μ_A plays a crucial role. This observation is due to Vilenkin [13].

Indeed, the axial current is not conserved because of the anomaly:

$$\partial_\mu\widehat{j}^\mu_A = \frac{e^2}{2\pi^2}E_\mu B^\mu = \frac{e^2}{4\pi^2}\partial_\mu\left(\epsilon^{\mu\lambda\rho\sigma}A_\lambda\partial_\rho A_\sigma\right) \tag{4}$$

so one may wonder whether a global thermodynamic equilibrium analysis holds. Indeed, this may happen under two circumstances:

- if the electromagnetic field is external and $E \cdot B = 0$

	B	μ_B	μ_A
C	■	■	✓
P	✓	✓	■
T	■	✓	✓
Rotation	■	✓	✓

Table 1. Symmetries of the density operator (3) which are maintained (checkmark) or broken (black square) by the parameters B, μ_B, μ_A.

- if the electromagnetic field is dynamical, provided that the axial current and the axial charge include the contribution of the Chern-Simons current reported on the right-hand-side of eq. (4).

The second option did not receive much attention in literature, except in ref. [5] but one should never forget that the dynamical electromagnetic field is always present. In fact, it can be shown that, in the limit of free electromagnetic field, the right hand side of the (4) corresponds to the difference between the number of right-handed and left-handed photons, so that one can loosely say that the total helicity of the fermion-boson system is overall conserved [14]. The possibility to study the Chiral Magnetic Effect in an equilibrium quantum statistical framework is very convenient in many respects. This feature, which can be seen as an expression tantamount to "the CME current is non-dissipative", makes the calculation of many quantities easier than, for instance, dissipative transport coefficients.

It is worth stressing that, even if in this formalism the anomaly is somewhat "hidden", it plays a crucial role to generate the chirality imbalance which is needed to give rise, after relaxation, to a finite μ_A. Because of the non-conservation of the axial current, a chiral charge can be created in a situation where, for some time, an external electromagnetic field is such that $E \cdot B \neq 0$, like in condensed matter experiments [15] or in the so-called magnetic reconnection [16]. In the QGP, the chiral charge is believed to arise from a topological non-Abelian sphaleron transition, and quantitative theoretical predictions are currently being provided [17].

Finally, I would like to make a point about the proper definition of anomalous effects. Indeed, it has been somewhat customary in literature to qualify as "anomalous" many non-dissipative currents which are not expected to appear in the traditional global thermodynamic equilibrium. For instance, the mean value of the axial current of the free Dirac field

$$j_A^\mu = \left(\frac{T^2}{6} + \frac{\mu^2}{2\pi^2} \right) \omega^\mu \tag{5}$$

in presence of vorticity, is qualified as "anomalous" in literature and associated to gravitational anomaly. In fact, this relation is not of anomalous origin because it can arise in *any* field theory involving Dirac fields, simply because it is allowed by the symmetry of the density operator with rotation:

$$\widehat{\rho} = \frac{1}{Z} \exp[-\widehat{H}/T + \omega \widehat{J}_z/T]$$

which was noted, again, for the first time by Vilenkin [18] and recently discussed in [19]. In my view, the qualification of anomalous is appropriate only when a current vanishes if the axial chemical potential μ_A does. The reason is explained above: without an anomalous divergence, it would be impossible to drive a chirality imbalance with external fields, that is to create a $\mu_A \neq 0$.

3. Polarization

Particles produced in relativistic heavy ion collisions are expected to be polarized in peripheral collisions because of angular momentum conservation. At finite impact parameter, the QGP has a finite angular momentum perpendicular to the reaction plane and some fraction thereof may be converted into spin of final state hadrons. Therefore, measured particles may show a finite mean *global* polarization along the angular momentum direction. In a fluid at local thermodynamic equilibrium, the polarization can be calculated by using the principle of quantum statistical mechanics, that is assuming that the spin degrees of freedom are at local thermodynamical equilibrium at the hadronization stage, much the same way as the momentum degrees of freedom.

The crucial role in the calculation of the polarization for the fluid produced in relativistic heavy ion collisions is played by the density operator. For a system at Local Thermodynamic Equilibrium (LTE), this reads [20]:

$$\widehat{\rho}_{\rm LE} = (1/Z) \exp\left[- \int_\Sigma {\rm d}\Sigma_\mu \left(\widehat{T}^{\mu\nu} \beta_\nu - \zeta \widehat{j}^\mu \right) \right] \tag{6}$$

where $\beta = (1/T)u$ is the four-temperature vector, $\widehat{T}$ the stress-energy tensor, $\widehat{j}$ a conserved current - like the baryon number - and $\zeta = \mu/T$. The mean value of a local operator $\widehat{O}(x)$ (such as, for instance the stress-energy tensor $\widehat{T}$, or the current $\widehat{j}$) at LTE:

$$O(x) = {\rm tr}(\widehat{\rho}_{\rm LE} \widehat{O}(x)) \tag{7}$$

and if the fields β,ζ vary significantly over a distance which is much larger than the typical microscopic length (indeed the *hydrodynamic limit*), then they can be Taylor expanded in the density operator starting from the point x where the mean value $O(x)$ is to be calculated. The leading terms in the exponent of (6) then become [20]:

$$\widehat{\rho}_{\rm LE} \simeq \frac{1}{Z_{\rm LE}} \exp\left[-\beta_\nu(x) \widehat{P}^\nu + \xi(x) \widehat{Q} - \frac{1}{4} (\partial_\nu \beta_\lambda(x) - \partial_\lambda \beta_\nu(x)) \widehat{J}^{\lambda\nu}_x + \frac{1}{2} (\partial_\nu \beta_\lambda(x) + \partial_\lambda \beta_\nu(x)) \widehat{L}^{\lambda\nu}_x + \nabla_\lambda \xi(x) \widehat{d}^\lambda_x \right]. \tag{8}$$

where the last two terms with the shear tensor and the gradient of ζ are dissipative and vanish at equilibrium. The ∇_λ operator stands for:

$$\nabla_\lambda = \partial_\lambda - u_\lambda u \cdot \partial$$

as usual in relativistic hydrodynamics. The term which is responsible for a non-vanishing polarization is the one involving the angular momentum-boosts operators $\widehat{J}_x$.

The polarization of particles in a fluid at LTE can in principle be obtained by calculating matrices like:

$$W_{\sigma,\sigma'} = {\rm tr}(\widehat{\rho}_{\rm LE} a^\dagger(p)_\sigma a(p')_{\sigma'})$$

where $a(p)_\sigma$ are the destruction operators of final state particles of four-momentum p and σ is the spin state index. Nevertheless, the exact calculation of W is a difficult one even with the expansion of $\widehat{\rho}_{\rm LE}$ and the mean polarization was obtained in ref. [21] by means of a different method, involving the spin tensor and an *ansatz* about the form of the covariant Wigner function at LTE (see also [22]). As a result, the mean spin vector of 1/2 particles with four-momentum p, turns out to be:

$$S^\mu(p) = -\frac{1}{8m} \epsilon^{\mu\rho\sigma\tau} p_\tau \frac{\int d\Sigma_\lambda p^\lambda n_F (1 - n_F) \varpi_{\rho\sigma}}{\int d\Sigma_\lambda p^\lambda f(x,p)} \tag{9}$$

where $n_F = (1 + \exp[\beta(x) \cdot p - \mu(x) Q/T(x)] + 1)^{-1}$ is the Fermi-Dirac distribution and $\varpi(x)$ is the *thermal vorticity*, that is:

$$\varpi_{\mu\nu} = -\frac{1}{2} \left(\partial_\mu \beta_\nu - \partial_\nu \beta_\mu \right) \tag{10}$$

The eq. (9) has been used in all numerical calculations of polarization, either based on the hydrodynamic model [23] or other approaches [24] and a good agreement with the data is observed. A crucial feature of the (9), and more in general of this effect, is that it predicts an almost equal polarization of particles and anti-particles (if quantum statistics are not important) for it is a statistical thermodynamic effect driven by local equilibration and not by an external C-odd field like the electromagnetic field. This distinctive feature is confirmed - modulo small deviations - by all measurements [1, 2, 3].

While an exact derivation of the formula (9) is still missing, this formula is the correct first-order expression in thermal vorticity in the non-relativistic limit and in the limit where quantum statistics can be neglected [21]. At a glance, it appears to be the most reasonable expression of the mean spin vector: it is linear in thermal vorticity, the spin is orthogonal to the four-momentum and it also has the right limit ($S^\mu = 0$) for the fully degenerate Fermi gas, i.e. $n_F = 1$.

To gain insight into the physics of polarization in a relativistic fluid, it is very useful to decompose the gradients of the four-temperature vector in the eq. (9). We start off with the seperation of the gradients of the comoving temperature and four-velocity field:

$$\partial_\mu \beta_\nu = \partial_\mu \left(\frac{1}{T}\right) + \frac{1}{T}\partial_\mu u_\nu$$

Then, we can introduce the acceleration and the vorticity vector ω^μ with the usual definitions:

$$A^\mu = u \cdot \partial u^\mu \qquad\qquad \omega^\mu = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}\partial_\nu u_\rho u_\sigma$$

The antisymmetric part of the tensor $\partial_\mu u_\nu$ can then be expressed as a function of A and ω:

$$\frac{1}{2}\left(\partial_\nu u_\mu - \partial_\mu u_\nu\right) = \frac{1}{2}\left(A_\mu u_\nu - A_\nu u_\mu\right) + \epsilon_{\mu\nu\rho\sigma}\omega^\rho u^\sigma$$

therafter plugged into the (9) to give:

$$\begin{aligned} S^\mu(x,p) &= \frac{1}{8m}(1-n_F)\epsilon^{\mu\nu\rho\sigma}p_\sigma \nabla_\nu(1/T)u_\rho - \frac{1}{8m}(1-n_F)\frac{1}{T}\epsilon^{\mu\nu\rho\sigma}p_\sigma A_\nu u_\rho \\ &+ \frac{1}{8m}(1-n_F)\,2\,\frac{\omega^\mu u\cdot p - u^\mu \omega\cdot p}{T} \end{aligned} \tag{11}$$

Hence, polarization stems from three contributions: a term proportional to the gradient of temperature, a term proportional to the vorticity ω, and a term proportional to the acceleration. Further insight into the nature of these terms can be gained by choosing the particle rest frame, where $p = (m, \mathbf{0})$ and restoring the natural units. The eq. (11) then certifies that the spin in the rest frame is proportional to the following combination (the natural constants are purposely restored):

$$\mathbf{S}^*(x,p) \propto \frac{\hbar}{KT^2}\gamma\mathbf{v}\times\nabla T + \frac{\hbar}{KT}\gamma(\boldsymbol{\omega} - (\boldsymbol{\omega}\cdot\mathbf{v})\mathbf{v}/c^2) + \frac{\hbar}{KT}\gamma\mathbf{A}\times\mathbf{v}/c^2 \tag{12}$$

where $\gamma = 1/\sqrt{1-v^2/c^2}$ and all three-vectors, including vorticity, acceleration and velocity, are observed in the particle rest frame.

The three independent contributions are now well discernible in eq. (12). The second term scales like $\hbar\omega/KT$ and is the one already known from non-relativistic physics, proportional to the vorticity vector seen by the particle in its motion amid the fluid, with an additional term vanishing in the non-relativistic limit. The third term is a purely relativistic one and scales like $\hbar A/KTc^2$; it is usually overwhelmingly suppressed, except in heavy ion collisions where the acceleration of the plasma is huge ($A \sim 10^{30}g$ at the outset of hydrodynamical stage). The first term, instead, is a new non-relativistic term [21] and applies to situations where the velocity field is not parallel to the temperature gradient. For ideal uncharged (thus relativistic) fluids, this term is related to the acceleration term because the equations of motion reduce to:

$$\nabla_\mu T = TA_\mu/c^2$$

Therefore, being the QGP a quasi-ideal fluid and almost uncharged at very high energy, the first and third term are tightly related. It can be shown that they contribute non-trivially to the final predicted polarization [25].

3.1. Open theoretical issues

As spin physics in relativistic fluids is a completely new subject, much work is still to be done and several theoretical questions are still to be answered. For instance, as has been mentioned, the formula (9) has been obtained by means of an educated *ansatz* of the Wigner function of the Dirac field and an exact derivation in the framework of quantum field theory is still to be found, even at the global thermodynamic equilibium, when the density operator reads:

$$\rho = \frac{1}{Z} \exp\left[-b_\mu \widehat{P}^\mu + \frac{1}{2}\varpi_{\mu\nu}\widehat{J}^{\mu\nu} + \zeta\widehat{Q}\right] \tag{13}$$

A major theoretical progress in this topic is compelling not just because of our aim at perfection, but in view of the by now demonstrated ability of the experiments to make differential measurements of polarization. Specifically, new results presented in this conference [3] have shown a disagreement with the predictions. Particularly, the azimuthal dependence of the mean spin component along the angular momentum S_J is at odds with the prediction of the hydrodynamic model [26] (see also [27] and references therein). Furthermore, it has been predicted, based again on the hydrodynamic model, that the component of the mean spin along the beam axis S_b, features oscillations in the azimuthal plane similar to the elliptic flow [28]. This pattern has been confirmed in a subsequent calculation based on AMPT [29]. Indeed, these oscillations, with the predicted periodicity, have been preliminarly observed by STAR and reported in this conference [3], yet with a flipped sign compared to the predictions. Needless to say, this finding would be a spectacular confirmation of the model if the sign matched the expectation.

Identifying the reason of the aforementioned discrepancies requires a deep theoretical investigation. The first possibility could be an incorrect choice of the initial hydro conditions, particularly the initial longitudinal flow profile. This hypothesis can be tested numerically with a dedicated study with presently available codes. A second possibility is the effect of decays of higher lying states. Indeed, all theoretical predictions of polarization as a function of momentum only include primary particles, that is Λ hyperons emitted from the hadronizing hypersurface. However, most Λ's stem from the decays of higher lying states and the polarization transfer is a non-trivial function of the momentum of the decay product. While global polarization transfer in decays producing a Λ hyperon has been calculated [30] and its impact evaluated numerically [28], nothing is known about it differentially in momentum space. Even if unlikely at first glance, a sign flip between primary and secondary Λ's is possible. A third possibility is that local equilibrium of spin degrees of freedom is not achieved as quickly as momentum and, consequently, the evolution of spin density is not just driven by thermal vorticity. Thereby, in order to describe a polarized relativistic, C-invariant fluid (such as the QGP at very high energy) a spin tensor $\mathcal{S}$ is necessary [31].

The latter is an intriguing option as it implies a profound consequence. The need of a spin tensor entails that different quantum stress-energy tensors are not equivalent to describe the fluid [32] and, more importantly, there should be measurable effects in the predicted polarization pattern. The hydrodynamic equations of a relativistic polarizeable fluid with particles and antiparticle would be supplemented by the angular momentum continuity equation:

$$\partial_\lambda \mathcal{S}^{\lambda,\mu\nu} = T^{\nu\mu} - T^{\mu\nu} \tag{14}$$

and that the hydrodynamic variables include also an antisymmetric tensor $\Omega_{\mu\nu}$ [31] playing the role of the thermodynamically conjugate variable of the spin tensor, much the same way as the chemical potential is the thermodynamically conjugate variable to a conserved charged current.

It is worth stressing that, with regard to polarization, the case of a non-relativistic fluid is in essence different from that of a non-relativistic fluid. In the former, polarization can be described by magnetization because of the full matter-antimatter asymmetry. In the latter, this is not enough and if particles and antiparticles are both polarized in the same direction, magnetization, which is C-odd, cannot make the job; a C-even polarization tensor, the spin tensor, is needed.

The transition from the familiar relativistic hydrodynamics to relativistic hydrodynamics with spin [33] is a major step forward and would certainly require a great amount of theoretical and numerical work. QGP offers the unique opportunity to test relativistic theories in the laboratory.

4. Conclusions

Polarization and Chirality have opened a new window in heavy ion physics. From a theory viewpoint, a new outlook, forcing us to rethink the foundations of relativistic hydrodynamics and kinetic theory in a fully quantum framework. The study of the quantum features of Quark Gluon Plasma have exciting connections with fundamental physics problems even beyond QCD.

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