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Directed flow of charm quarks as a witness of the initial strong magnetic field in ultra-relativistic heavy ion collisions



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ABSTRACT

Ultra-relativistic Heavy-Ion Collision (HIC) generates very strong initial magnetic field (\vec{B}) inducing a vorticity in the reaction plane. The high \vec{B} influences the evolution dynamics that is opposed by the large Faraday current due to electric field generated by the time varying \vec{B} . We show that the resultant effects entail a significantly large directed flow (v_1) of charm quarks (CQs) compared to light quarks due to a combination of several favorable conditions for CQs, mainly: (i) unlike light quarks formation time scale of CQs, $\tau_f \simeq 0.1$ fm/c is comparable to the time scale when \vec{B} attains its maximum value and (ii) the kinetic relaxation time of CQs is similar to the QGP lifetime, this helps the CQ to retain the initial kick picked up from the electromagnetic field in the transverse direction. The effect is also odd under charge exchange allowing to distinguish it from the vorticity of the bulk matter due to the initial angular momentum conservation; conjointly thanks to its mass, $M_c >> \Lambda_{QCD}$, there should be no mixing with the chiral magnetic field produced in HIC.

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The properties of the hot and dense phase of matter referred to as Quark–Gluon Plasma (QGP) expected to be produced in nuclear collisions at relativistic energies are governed by light quarks and gluons [1,2]. The heavy flavors namely, charm and bottom quarks play crucial roles in probing the QGP [3]. The special role of CQs as a probe of the QGP resides on the fact that their mass, *M* is significantly larger than the typical temperatures (*T*) achieved in QGP and the QCD scale parameter (Λ_{QCD}) *i.e.* $M \gg T$, Λ_{QCD} , therefore, the production of heavy quarks is essentially limited to the primordial stage of a heavy-ion collision with a formation time $\tau_c \sim 1/2 M_c \simeq 0.08$ fm/c. In such a scenario the probability of CQs getting annihilated or created during the evolution is much smaller compared to light quarks and gluons. As a consequence CQs witness the entire space–time evolution of the system and can act as an effective probe of the created matter.

The two main experimental observables related to CQs which have been extensively used as QGP probes are: (i) the nuclear suppression factor, R_{AA} which is the ratio of the p_T spectra of heavy flavored hadrons (D and B) produced in nucleus + nucleus collisions to those produced in proton + proton collisions appropriately scaled at a given $\sqrt{s_{NN}}$ and (ii) the elliptic flow, $v_2 = \langle \cos(2\phi_p) \rangle$, a measure of the anisotropy in the angular distribution that corresponds to the anisotropic emission of particles with respect to the reaction plane. In the present work we demonstrate that the directed flow $v_1 = \langle \cos(\phi_p) \rangle = \langle p_x/p_T \rangle$ of CQs is a superior probe to estimate the magnetic field generated in non-central HICs.

Several theoretical efforts have been made within the ambit of Fokker Planck [4–18] and relativistic Boltzmann transport approaches [19–27] to calculate R_{AA} [28–31] and v_2 [30]. Essentially all the models show some difficulties in simultaneously describing the $R_{AA}(p_T)$ and $v_2(p_T)$ and such a trait is not only present at RHIC but also appears in a stronger way at LHC energy. However it has been shown in Ref. [32] that a nearly constant drag or slightly rising with T in the range of $\Gamma \sim 0.15-0.2$ fm⁻¹ is able to simultaneously describe $R_{AA}(p_T)$ and $v_2(p_T)$ at least at RHIC energy, while at LHC energy it still remains uncertain also due to the large experimental error bars (see [6,7,33]).

In recent years it has been recognized that a very strong magnetic field is created at early times in heavy ion collisions [34, 35]. The impact of the magnetic field was explored mainly in re-

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lation to the chiral magnetic effect [36,37], but also to jet energy loss [38], to J/Ψ elliptic flow [39] as well as to thermal photon and dilepton productions [40], and very recently to the CQ diffusion coefficients [41,42]. The estimated values of the initial field strengths, $eB_y\sim 5\,m_\pi^2\,$ and $\sim 50\,m_\pi^2\,$ at RHIC and LHC energies respectively, which is several orders of magnitude higher than the values predicted at the surface of magnetars. Since the CQs are produced at the early stage of HICs, we argue that their dynamics will be affected by such a strong magnetic field and they will be able to retain these effects till its detection as D mesons in experiments. The \vec{B} field generated in non-central HICs is dominated by the \vec{v} -component which induces a Faraday current in the xz plane. In particular due to the expansion along the z axis the Lorentz force is directed along the negative (positive) \vec{x} direction in the forward (backward) rapidity region for positively (negatively) charged quarks. This can be seen as a classical Hall effect that generates a directed transverse flow. In addition to the Hall effect the time dependence of \vec{B} generates a large electric field due to Faraday effect according to $\vec{\nabla} \times \vec{E} = -\partial \vec{B} / \partial t$. The induced Faraday current opposes the drift due to the magnetic field. The combination of the two effects leads to a finite v_1 .

To study quantitatively the global dynamics of CQs we solve the relativistic Langevin equation in an expanding QGP background. In this work the initial conditions for solving the Langevin equation for the CQs and the relativistic transport code for the background are constrained by the experimental data on the $R_{AA}(p_T)$ and v_2 of D meson [32] and the transverse momentum spectra and the elliptic flow of the bulk, see [43-46] for more details. The dynamics of the CQs in the QGP is largely determined by the drag coefficient, Γ . We have set a weak T dependence to Γ in the interval 0.15–0.2 fm⁻¹ which helps in reproducing the $R_{AA}(p_T)$ and $v_2(p_T)$ for D meson at RHIC and LHC collisions reasonably well [32]. The initial condition for the bulk in the transverse r-space is taken from the Glauber model assuming boost invariance along the longitudinal direction. The value of the maximum temperature in the center of the fireball at the initial time, $\tau_0 = 0.2$ fm/c is set as $T_0 = 580$ MeV. The initial spatial and momentum distributions of the CQs are set respectively by the N_{col} and the FONLL scheme of the charm production in proton + proton collisions [47,48] at the same $\sqrt{s_{\rm NN}}$.

Intensive studies have been performed in the recent years to determine the electromagnetic field generated in ultra-relativistic HIC [40,49,50]. In the present work we aim to make the first study on the impact of the electromagnetic field on heavy quark dynamics under the assumption of a constant electrical conductivity (σ_{el}) of the QGP. This will enable us to – obtain analytic solutions, highlight the core physics and avoid further numerical complications. We refer to [51] for the details of the space-time dependent solutions of \vec{E} and \vec{B} (see also [40,50]). With the *z* and *x* axes along the beam and impact parameter directions respectively the \vec{B} generated in non-central HIC will point along the *y*-axis on the average, *i.e.* $\mathbf{B} = B_y \hat{e}_y$ while the other components, $B_x = B_z = 0$. The electric field, **E** due to Faraday effect will align along the *x* axis.

The density $\rho_{\pm}(\vec{x_{\perp}})$ of the protons at $\vec{x_{\perp}}$ in the transverse plane can be estimated by projecting the probability distribution of the protons homogeneously distributed in a spherical nucleus moving either in the +z or -z. In a collision with impact parameter $b \neq 0$ the + and - signs indicate the spectators moving along + and z directions respectively and the total magnetic field is the sum of $B_{\nu}^{+,-}$ generated by Z point-like charges as [36]

$$eB_y = -Z \int_{-\frac{\pi}{2}}^{\frac{\pi}{2}} d\phi' \int_{x_{\rm in}(\phi')}^{x_{\rm out}(\phi')} dx'_{\perp} x'_{\perp} \rho_{-}(x'_{\perp})$$

$$\times \left(eB_{\nu}^{+}(\tau,\eta,x_{\perp},\phi) + eB_{\nu}^{-}(\tau,\eta,x_{\perp},\phi)\right),\tag{1}$$

where ϕ is the azimuthal angle, $\tau \equiv \sqrt{t^2 - z^2}$ is the proper time, $\eta \equiv \arctan(z/t)$ is the space–time rapidity, and $x_{\rm in}$ and $x_{\rm out}$ are the endpoints of the x'_{\perp} integration regions that define the crescent-shaped loci where one finds protons either moving + or -z directions but not both. These are given by

$$x_{\rm in/out}(\phi') = \pm \frac{b}{2}\cos(\phi') + \sqrt{R^2 - \frac{b^2}{4}\sin^2(\phi')},$$
 (2)

where *R* is the radius of the nucleus, and *b* is the impact parameter of the collision. The main ingredient of Eq. (1) is the magnetic field B_y at an arbitrary space-time point $(t, \vec{x_\perp}, \eta)$ generated by a single charge located at x'_{\perp} moving in the +(-)z direction with velocity $\vec{\beta}$. The B_y^+ can be written as,

$$eB_{y}^{+}(\tau,\eta,x_{\perp},\phi) = \alpha \sinh(Y_{b})(x_{\perp}\cos\phi - x'_{\perp}\cos\phi')$$
$$\frac{\sigma_{el}\frac{|\sinh(Y_{b})|}{2}\xi^{\frac{1}{2}} + 1}{\xi^{\frac{3}{2}}}e^{A}, \qquad (3)$$

where $\alpha = e^2/(4\pi)$ is the electromagnetic coupling, $Y_b \equiv \arctan(\beta)$ is the beam rapidity of the + mover, *A* and ξ stand for

$$A \equiv \frac{\sigma_{el}}{2} \left(\tau \sinh(Y_b) \sinh(Y_b - \eta) - |\sinh(Y_b)| \xi^{\frac{1}{2}} \right)$$
(4)

$$\xi \equiv \tau^2 \sinh^2(Y_b - \eta) + x_\perp^2 + x_\perp'^2 - 2x_\perp x_\perp' \cos(\phi - \phi').$$
(5)

We note that the time evolution of the magnetic field is determined by σ_{el} which drives the magnitude of *A* in Eq. (3). The following calculations are performed for $\sigma_{el} = 0.023 \text{ fm}^{-1}$ obtained from lattice QCD calculations [52–54] in the temperature range around $\sim 2 T_c$.

In a similar approach the *x*-component of the electric field produced by the charges moving along + z direction can be obtained as:

$$eE_{x}^{+}(\tau,\eta,x_{\perp},\phi) = eB_{y}^{+}(\tau,\eta,x_{\perp},\phi)\operatorname{coth}(Y_{b}-\eta), \qquad (6)$$

which again has to be convoluted with the transverse charge distribution $\rho_{\pm}(x_{\pm})$ as for the magnetic field. The other components of the electromagnetic field averaged over initial conditions will vanish or become quite small as in the case of the x-component of the electric field E_z in the region between the two colliding beams where the plasma is formed. However, it is has been shown in [49] that the large fluctuations in event by event central collisions (not of interest here) can generate other components of the fields with magnitudes comparable but generally smaller than to B_{ν} and E_{x} . Also the positive (negative) direction of the B_{y} -field here is conventional and in the experiments an event-by-event analysis has to be done to find a non-vanishing flow. Furthermore, in principle one should also include the electromagnetic field generated by the participant protons, however, it has been shown in [51] that its magnitude is sub dominant especially in the initial stage that plays the leading role for the directed flow considered here.

In Fig. 1 (upper panel) we display the time evolution of the magnetic field $\mathbf{B} = B_y \hat{e}_y$ at $\vec{x_\perp} = 0$ for various η for Pb + Pb at $\sqrt{s_{\rm NN}} = 2.76$ TeV for b = 9.5 fm with $\sigma_{el} = 0.023$ fm⁻¹. The electric field vanishes at this position due to symmetry. An important factor for sizable directed flow is certainly the fact that CQs are produced in the early stage. In fact in this region the B_y reduces by about an order of magnitude between t = 0.1 fm/c to t = 1 fm/c. In Fig. 1 (lower panel) we depict the time dependence of both B_y (black line) and electric field $\mathbf{E} = E_x \hat{e}_x$ (red line) at $\mathbf{x}_\perp = 0$ and



Fig. 1. (Color online.) Time dependence of B_y and E_x fields for $\sigma_{el} = 0.023 \text{ fm}^{-1}$ in Pb + Pb collisions at $\sqrt{s_{NN}} = 2.76$ TeV for b = 9.5 fm at $x_{\perp} = 0$. Upper panel: eB_y for different space rapidities η ; Lower panel: time evolution of both the magnetic field eB_y (black) and the electric field eE_x (red) at forward rapidity $\eta = 1.0$.

 $\eta = 1.0$. We note that for t < 1 fm/c there is a large difference between the B_y and E_x , although they become equal at later time. We will see that this plays an important role in determining the sign and the size of v_1 .

The dynamics of the CQs, with charge q and momentum p, is governed by the Langevin equation in the presence of electromagnetic field, given by

$$\dot{\mathbf{x}}\left(t\right) = \frac{\mathbf{p}}{E} \tag{7}$$

$$\dot{\mathbf{p}}(t) = -\Gamma \mathbf{p}(t) + \mathbf{F}(t) + \mathbf{F}_{ext}(t), \qquad (8)$$

where the first term represents the dissipative force and the second term represents the fluctuating force *F*(*t*) regulated by the diffusion coefficient *D*. The third term in Eq. (8) represents the external Lorentz force due to the electric and magnetic fields. We study the evolution with a standard white noise ansatz for *F*(*t*), *i.e* $\langle \mathbf{F}(t) \rangle = 0$ and $\langle \mathbf{F}(t) \mathbf{F}(t') \rangle = D\delta(t - t')$. The ensemble $\langle ... \rangle$ denotes the averaging of many trajectories for *p* each consisting of different realizations of *F* at each time step. To solve the Langevin equation for an expanding system one needs to move to the local rest frame of the background fluid [3,5], where an element moving with velocity **v** with respect to the laboratory frame will be subjected to both **E**' and **B**' as determined by Lorentz transformations. The *F*_{ext} in the fluid rest frame will be

$$\mathbf{F}_{ext} = q\mathbf{E}' + \frac{q}{E_p} \left(\mathbf{p} \times \mathbf{B}' \right) \tag{9}$$

where $E_p = \sqrt{p^2 + M^2}$ is the energy of the heavy quark with momentum *p*.

In Fig. 2 we show the resulting directed flow v_1 as a function of the rapidity of charm black (solid line) and anti-charm quarks



Fig. 2. (Color online.) Directed flow v_1 as a function of the rapidity in Pb + Pb collision at $\sqrt{s_{NN}} = 2.76$ TeV for b = 9.5 fm for D meson $[c\bar{q}]$ at $p_T > 1$ GeV black (solid) and anti-D meson $[\bar{c}q]$ (dashed) line at $t = t_{f.o.}$ (see text). Red dash-dot (blue dash-dot-dot) line indicates v_1 of D meson at t = 2 fm/c (t = 5 fm/c).

(dashed line). We can see that there is a substantial v_1 at finite rapidity with a peak at $y \simeq 1.75$. The flow is negative for positive charged particle (charm) at forward rapidity which means that the Hall drift induced by the magnetic field dominates over the displacement caused by the Faraday current associated with the time dependence of the magnetic field. This is a non-trivial result and it is partially due to the fact that the formation time of CQs is very close to the time at which the magnetic field attains its maximum, causing a large drift due to Hall effect.

We have followed the dynamics up to t = 12 fm/c, but observed that the directed flow saturates already at $t \simeq 1-2$ fm/c for |y| < 0.5-1, and at $t \simeq 5$ fm/c for |y| < 1.5 (Fig. 2). Therefore, the slope $dv_1/dy|_{y=0} \simeq -1.75 \cdot 10^{-2}$ is determined in the very early stage of the collision $t \leq 1-2$ fm/c. The time scale of the saturation of v_1 as a function of y follows the persistence of the B_y and E_x fields with increasing η shown in Fig. 1 (upper panel). Therefore, the very early stage and its sign and magnitude is essentially controlled by the large value of B_y for $t \leq 1.0$ fm/c. The predicted value of $v_1(y)$ for D meson $[c\bar{q}]$ is quite large and the odd behavior of D/D doubles the effect that can be measured. Also it would be a distinctive signal of the electromagnetic field, distinguishable from the $v_1(y)$ that can be generated by angular momentum conservation as studied in [55].

It is important to stress that a main feature of the CQs that turns out to favor the formation of a sizable directed flow is the relative large equilibration time w.r.t. light quarks. In fact, the relaxation time of CQs can be estimated as $\tau_c^{eq} \simeq 1/\Gamma \approx 5-8$ fm/c which is much larger than the light quark and gluon equilibration time, $\tau_{QGP}^{eq} \approx 0.5-1$ fm/c.

In Fig. 3 we show a study of the strong dependence of the transverse flow on the interaction strength given by the drag coefficient Γ and plotted in term of the equilibration time defined as $\tau_{eq} = 1/\Gamma$.

The strong dependence of v_1 on Γ is evident from the variation of $|dv_{1c}/dy|$ with τ_{eq} as displayed in Fig. 3. The quantity, $|dv_{1c}/dy|$ for CQ with $\tau_c^{eq} \simeq 1/\Gamma \approx 5-8$ fm/c is at least two orders of magnitude larger than the corresponding value for light quarks with $\tau_{eq} \sim 0.6$ fm/c [51]. This is due to the fact that the transverse kick exerted by the electromagnetic field during the time interval, $\tau_{e.m.}$ on the thermalized light quarks (unlike CQ which is out-of-equilibrium) is damped by its random interaction in the medium with similar durability. However, the lowest points in Fig. 3 may not be taken as a realistic estimate for v_1 of light quarks, because for that we have to keep in mind at least three other aspects: the dynamics of light quarks cannot be appropriately studied by us-



Fig. 3. (Color online.) Absolute value of the slope of the charm transverse flow $|dv_1/dy|$ around mid-rapidity for Pb + Pb at $\sqrt{s_{NN}} = 2.76$ TeV for b = 9.5 fm as a function of the inverse of the drag coefficient Γ and for two different values of the thermalization time $\tau_0 = 0.2$ fm/c (circles) and $\tau_0 = 0.6$ fm/c (squares).

ing Langevin dynamics as is done usually for heavy quarks, the light hadrons originate abundantly also from the hadronization of gluons which are not directly affected by the electromagnetic interaction and their initial momentum distributions is quite different from that of CQs. All these aspects cause a further significant reduction in the transverse flow of light hadrons which will be discussed in a future work.

In this context it is also important to mention that initially the charm quarks are in a non-equilibrium stage and due to the scatterings with the medium experience a significant acceleration in the first fm/c. Such an acceleration is proportional to the drag coefficient and it can reach values similar to the longitudinal expansion rate $1/\tau$ at τ_0 for values of the drag corresponding to $\tau_{eq} = 1/\Gamma \simeq 0.5$ fm/c. From Fig. 3 we see that in such a case the strong longitudinal acceleration implies a nearly complete damping of the transverse kick that would be induced by the magnetic field.

A last aspect we want to point out is that certainly the strength of the electromagnetic field is important to have a sizeable transverse flow, but the underlying dynamics is more subtle. In Fig. 4 we display the $v_1(y)$ that is generated if we switch off artificially the electric field and keep the action of the magnetic field on (Hall drift only). We notice that in such a situation the $v_1(y)$ (black lines) generated is much larger than the one displayed in Fig. 2. We observe also that when only the electric field is considered the effect of the Faraday current generates v_1 with opposite sign but a magnitude similar to the Hall drift. The $v_1(y)$ in Fig. 2 even if not exactly equal to the difference between the Hall drift and the Faraday current calculated separately, as in Fig. 4, differs from it only by at most 5–10%. We understand that the value of v_1 is not only decided by the magnitude of the fields, but depends critically on the balance between \vec{E} and \vec{B} fields. In particular, the magnitude of the magnetic Hall drift, depending on the absolute magnitude of B_{ν} , is large at the formation time of the CQs. This entails a dominance of the Hall drift that is kept till the end of the evolution. Instead light quarks likely fail to feel the presence of the early high magnitude of $B_y > E_x$ due to their late formation. In fact, looking at Fig. 1 (lower panel) and Fig. 4 it is straightforward to envisage that if the charged particles would be produced at $t \simeq 1$ fm/c or if the simulation of the dynamics starts at similar times then the electric and magnetic field nearly compensate their effects, consequently $v_1(y)$ with smaller magnitude, see also Fig. 3.



Fig. 4. (Color online.) Black (Magenta) line shows the variation of $v_1(y)$ with y generated by the drift due to Hall effect ($B_y \neq 0$, $E_x = 0$) (generated by Faraday effect ($E_x \neq 0$, $B_y = 0$)) in Pb + Pb collision at $\sqrt{s_{NN}} = 2.76$ TeV for b = 9 fm for $p_T > 1$ GeV. The $v_1(y)$ for charm (anti-charm) is denoted by the solid (dashed) line.

One may also wonder what can be expected for bottom quarks. Granted they have a factor of 2 smaller coupling to the e.m field due to the charge $\pm 1/3$, the larger mass leads to a significant damping of the Lorentz Force proportional to p/E_p . A preliminary calculation shows that this determines a nearly exact balance between the Hall drift and Faraday current resulting in a v_1 that is about 4-5 times smaller the charm quark one, but its value critically depends on the details of the drag coefficient, initial time τ_0 and p_T distribution, that currently under scrutiny and will be presented in a future work.

We have also checked the impact of the electromagnetic field on $R_{AA}(p_T)$ and $v_2(p_T)$ and found that the former are not altered by the electromagnetic force; while an effect of the B-field on v_2 can come indirectly from the anisotropy induced in the bulk as conjectured in [41].

In summary, the present study suggests that v_1 of CQs can be considered as an efficient probe to characterize the evolving magnetic field produced in ultra-relativistic HIC. The time evolution of the field is determined by the electrical conductivity of QGP created in such collisions. Our central focus has been to show that the electromagnetic field can generate a sizable v_1 for CQs and hence for D meson, thanks to several concurring favorable effects for this to happen. The formation time of CQs is of about $au_{form} \sim 0.1~{
m fm/c}$ that is when the intensity of the \vec{B} field is maximum, even more important aspect is that the dynamics at time $t \leq 1.0$ fm/c is governed by the opposite action of E_x and B_y provides significant amount of net flow. Furthermore, the CQs, due to their large relaxation time in contrast to light quarks, are capable of retaining the memory of the initial non-equilibrium dynamics more effectively and hence lending stronger signal of the early magnetic dynamics. Furthermore, a large number of light hadrons originate from the gluons, not directly coupled to the electromagnetic field. Also in this respect CQ would provide a much cleaner and direct probe of the magnetic field dynamics. All these favorable conditions largely overwhelm the small suppression of the Lorentz force by a factor, p/E_p due to their finite mass. In addition, for $M_c >> \Lambda_{QCD}$ the CQ dynamics is not significantly affected by the chiral dynamics, therefore, the splitting between charge and anti-charge would not mix with the Chiral Magnetic Effect (CME) and/or with possible Chiral Vortex Effect (CVE) that can also generate a matter/antimatter splitting [56,57]. Thus, CQs would provide an independent way to scrutinize and quantify the initial magnetic field which can in turn also contribute to a more quantitative assessment of the CME and CVE.

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