

Asymmetric muon-antimuon emission from Z^0 decays: A potential magnetometer in relativistic heavy-ion collisions

Alejandro Ayala^{1,2,3}, Ana Julia Mizher^{2,3,4,5} and Javier Rendón¹

¹*Instituto de Ciencias Nucleares, Universidad Nacional Autónoma de México, Apartado Postal 70-543, CdMx 04510, Mexico*

²*Laboratório de Física Teórica e Computacional, Universidade Cidade de São Paulo, 01506-000 São Paulo, SP, Brazil*

³*Instituto de Física Teórica, Universidade Estadual Paulista, Rua Dr. Bento Teobaldo Ferraz, 271 Bloco II, 01140-070 São Paulo, SP, Brazil*

⁴*Centro de Ciencias Exactas and Departamento de Ciencias Básicas, Facultad de Ciencias, Universidad del Bío-Bío, Casilla 447, Chillán, Chile*

⁵*Instituto de Física, Universidade de SP, Rua do Matão, 1371, 05508-090 São Paulo, SP, Brazil*



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We show that a very clear signal of the presence of a strong magnetic field during the early stage of a high-energy heavy-ion collision is provided by the decay of the Z^0 into dimuon pairs. We find that the process is highly anisotropic, producing pairs mainly out of plane, as signaled by a negative value of v_2 , and leads to an antimuon transverse momentum distribution which peaks at a higher value of the transverse momentum compared to the peak of the muon transverse momentum distribution. We also show that the process does not produce a significant distortion of the Z^0 spectral function. The signal can be identified by comparing the dimuon-invariant mass and the individual muon and antimuon spectra produced in semicentral heavy-ion collisions with the corresponding scaled spectra produced in $p + p$ collisions at the Z^0 peak.

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I. INTRODUCTION

Heavy-ion collisions are an excellent tool for exploring the properties of hadron matter subject to extreme conditions. In recent times, it has been recognized that such conditions include not only large temperatures and densities but also the possible presence of strong electromagnetic fields. However, in spite of the many efforts aimed at identifying clear probes directly linked to the presence of such fields, it is fair to say that up to now the clearest signal has been provided only by the possible detection of the linear Breit-Wheeler process in ultraperipheral collisions [1–3]. For the kind of collisions that produce hot hadron matter, the impact of electromagnetic fields on the quark-gluon plasma (QGP) evolution was recently reported in terms of a possible correlation between the properties of the measured directed flow coefficients of charged pions, kaons, protons, and antiprotons and the strength of the magnetic field in semicentral collisions [4–6].

A putative magnetic field can, in principle, affect all stages of the evolution of a heavy-ion reaction. However, the field strength fades rather rapidly with time [7–11]. Most of the estimates of the field strength and its time evolution come from classical calculations, as opposed to true quantum calculations. Reference [12] suggests that the classical description

is valid and shows it by computing nucleon resonance production and its decay in ultraperipheral heavy-ion collisions from both a classical and a quantum description, finding that the results are consistent within a few percent. The fast decrease of the field strength with time makes it more difficult to identify its imprints with probes produced during the QGP or hadronic stage [13]. A better chance to identify the presence of such fields can be provided by penetrating probes produced during the earliest stages, where the field is close to its peak intensity. Good probes with these properties are the dilepton pairs that originate from the decay of Z^0 bosons [14,15]. Direct photons—when produced during preequilibrium when gluons are abundant—from magnetic-field-induced gluon fusion and splitting also serve as good penetrating probes carrying the information of the magnetic field strength [16–19]. A thermalized magnetized medium that exhibits magnetic fluctuations [20–23] has been shown to enhance the photon production rate by approximately one order of magnitude, while leaving the angular distribution of the emitted photons largely unchanged [24]. Another imprint of the presence of a strong magnetic field during the preequilibrium may be the short time required for the development of the hydrodynamical regime, given that a magnetic field helps to temper the originally large anisotropic pressures in the longitudinal and transverse directions with respect to the beam, as compared to the case without the field [25]. Other processes from where the strength of the magnetic field during the collision can be inferred include the emission of photons by bremsstrahlung and pair annihilation in a quark-gluon plasma [26–30]; production of electromagnetic radiation from the QED $\times$ QCD conformal anomaly [31]; fluctuations of the gluon field

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coupled to the photon stress tensor [32]; Cherenkov emission in a strong magnetic field [33] and from a quark-gluon plasma in a weak magnetic field coupled to the longitudinal dynamics in the background medium [34]; thermal dileptons emitted from a quark-gluon plasma in a weak field [35] and in a time-dependent magnetic field [36]. Signals of deconfinement from photon production in a long-lived chromomagnetic background, triggered by a flash of a strong electromagnetic field, were studied in Ref. [37]. Holographic methods have also been used to describe photon production from a strongly coupled plasma in strong magnetic fields [38–41]. A pedagogical review of various properties of electromagnetic fields and its connections with anomalous transport phenomena and experimental signatures in heavy-ion collisions is provided in [42]. A more recent review of the effects of strong electromagnetic fields on hadron matter is found in Ref. [43].

With the goal of providing a very clear signal of the presence of magnetic fields in semicentral heavy ion collisions, in this paper we study the production of dilepton pairs from Z^0 decays in the presence of a strong magnetic field. Dilepton production from virtual photons in a magnetized QGP and from ρ decays in a magnetized hadronic medium has been studied in Refs. [44,45] and [46,47], respectively. However, the novel scenario we present here consists of studying the production of Z^0 at the beginning of the reaction and its subsequent decay during the preequilibrium stage. The advantage of such a scenario is that, given the large Z^0 mass, these particles are produced at very early times by means of hard quark-antiquark annihilation, and rapidly decay into dilepton pairs at times where the field is still at its largest intensity. Because dileptons are a penetrating probe, they escape from the reaction, carrying a direct imprint of the field intensity during the early stages. We show that the imprint is not a distortion of the spectral density; despite the large magnitude of the magnetic field, the Z^0 mass and width do not suffer significant modifications, even for the largest intensities thought to be present during the reaction. This happens, as we show, because the Z^0 mass and width magnetic field-induced modifications are suppressed by a factor $(m_q/M_Z)^2$, where m_q is the quark mass and M_Z the Z^0 mass. However, the dilepton emission is enhanced by a factor of the field intensity squared. As we also show, the dilepton emission occurs mostly along the direction of the magnetic field, which can be quantified as a negative contribution to the v_2 of the dilepton azimuthal distribution. The emission rate is suppressed by a factor $(m_l/M_Z)^2$, where m_l is the lepton mass, and thus, instead of looking at dielectron pairs, we concentrate on dimuon emission. Given the large mass difference between the muon and the Z^0 at the Z^0 peak, muons behave almost as massless particles. This, together with the fact that they come from an electroweak decay process and move in a strong magnetic field, provides a correlation between their polarization, their direction of motion, and consequently their transverse (to the beam) momentum p_T . We show that antimuons are produced with a harder p_T distribution compared to muons. Both the anisotropic and the different p_T distributions of muons and antimuons provide clear signatures of the presence of the magnetic field, as well as a handle to infer the field strength

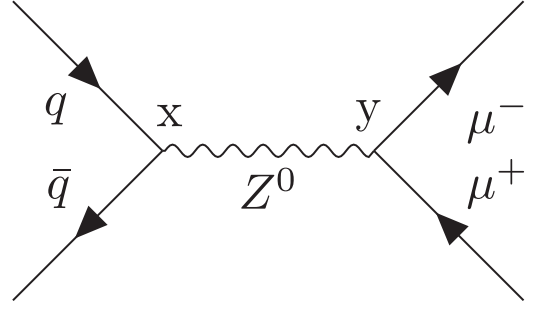


FIG. 1. Feynman diagram representing the current-current correlation function that describes the production of a Z^0 from quark-antiquark annihilation in the nuclear collision at a space-time point x followed by its subsequent decay into a dimuon pair at another space-time point y , in the presence of a strong magnetic field.

during the preequilibrium stage of the collision, which can be sought experimentally.

II. DIMUON PRODUCTION RATE

The starting point is the calculation of the current-current correlation function that describes the production of a Z^0 from an initial hadronic state (quark-antiquark annihilation in the nucleus-nucleus collision) at a given space-time point x followed by its subsequent decay into a dimuon pair in the final state at another space-time point y , in the presence of a strong magnetic field. This is depicted in Fig. 1. Integration over the space-time points x and y at the production and decay vertices provides the amplitude of the process. The probability of dimuon production, per unit volume and time, is obtained by squaring this amplitude and dividing it by the space-time volume where the reaction is active. This strategy has been successfully followed to compute dilepton production at finite temperature [48] and at finite temperature and density [49] from decays of ρ^0 through vector dominance.

The calculation can be written in terms of the product of a tensor describing the process on the hadron side, $H^{\mu\nu}$, and a tensor describing the process on the lepton side, $L_{\mu\nu}$. To compute the hadron tensor, we can resort to the optical theorem, which amounts to computing the imaginary part of the magnetic field modified Z^0 propagator, which in turn can be obtained from the Z^0 polarization tensor in the presence of the magnetic field, $\Pi^{\mu\nu}$. We work at one-loop order and in the lowest Landau-level (LLL) approximation for the quark propagators in the presence of a constant magnetic field pointing along the $\hat{z}$ direction. This is depicted in Fig. 2 and is the relevant scenario in the case the magnetic field is much larger than the quark mass squared. The Z^0 polarization tensor in the presence of a magnetic field of arbitrary strength was computed in Ref. [50]. In our approximation, the polarization tensor can be expressed as

$$\begin{aligned} \Pi^{\mu\nu}(q_{\parallel}^2, q_{\perp}^2) &= \Pi_{\parallel}(q_{\parallel}^2, q_{\perp}^2)P_{\parallel}^{\mu\nu} + \Pi_{\perp}(q_{\parallel}^2, q_{\perp}^2)P_{\perp}^{\mu\nu}, \\ P_{\parallel}^{\mu\nu} &\equiv g_{\parallel}^{\mu\nu} - q_{\parallel}^{\mu}q_{\parallel}^{\nu}/q_{\parallel}^2, \quad P_{\perp} \equiv q_{\parallel}^{\mu}q_{\parallel}^{\nu}/q_{\parallel}^2, \end{aligned} \quad (1)$$

where $g_{\parallel}^{\mu\nu} = \text{Diag}(1, 0, 0, -1)$ is the metric tensor in the parallel subspace and $q_{\parallel}^{\mu} = (q_0, 0, 0, q_3)$, $q_{\perp}^{\mu} = (0, q_1, q_2, 0)$

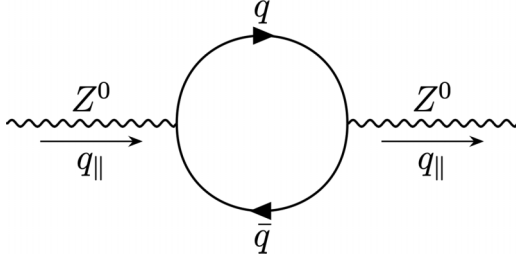


FIG. 2. Feynman diagram representing the one-loop Z^0 self-energy. In the strong field limit the quark propagators are taken in the lowest Landau level.

are the Z^0 longitudinal and transverse four-momentum vectors containing only the timelike and third component and the transverse components, respectively. The tensor structures on the right-hand side of Eq. (1) represent the possible polarization states for the Z^0 . Notice that the first (second) tensor structure is transverse (longitudinal) with respect to the four-momentum $q_\parallel^\mu$, and consequently the two possible polarization states can either be transverse or longitudinal with respect to the propagation four-momentum containing only the third component of the linear momentum. In a semi-classical picture, this means that the quark-antiquark pair that produced the Z^0 moved in a tightly bound orbit around a given field line. As a consequence, the transverse momentum components can be considered as averaging to zero, although the magnitude of the transverse momentum is not vanishing and is of order of the magnetic field strength. This is the well-known dimensional reduction introduced by working in the LLL approximation. The imaginary part of the functions $\Pi_\parallel$ and Π_L are given by

$$\text{Im}\{\Pi_\parallel, \Pi_L\} = g_Z^2 \frac{|e_q B|}{4\pi} e^{-\frac{q_\perp^2}{2|e_q B|}} \{-C_V^2, C_A^2\} F(q_\parallel^2) \theta(q_\parallel^2 - 4m_q^2),$$

$$F(q_\parallel^2) \equiv \left(\frac{m_q^2}{q_\parallel^2}\right) \sqrt{\frac{q_\parallel^2}{q_\parallel^2 - 4m_q^2}}, \quad (2)$$

where g_Z is the weak neutral coupling, e_q is the quark electric charge, and C_V and C_A are the Z^0 vector and axial-vector couplings to the quarks. Also, notice that Eq. (2) exhibits the usual energy threshold for pair production. The real parts of the functions $\Pi_\parallel$ and Π_L are given by

$$\text{Re } \Pi_\parallel = g_Z^2 \frac{|e_q B|}{8\pi^2} e^{-\frac{q_\perp^2}{2|e_q B|}} \{C_V^2 [F(q_\parallel^2)G(q_\parallel^2) - 1] - C_A^2\},$$

$$\text{Re } \Pi_L = -g_Z^2 \frac{|e_q B|}{8\pi^2} e^{-\frac{q_\perp^2}{2|e_q B|}} C_A^2 F(q_\parallel^2) G(q_\parallel^2),$$

$$G(q_\parallel^2) \equiv \ln \left(\frac{\sqrt{q_\parallel^2 - 4m_q^2} - \sqrt{q_\parallel^2}}{\sqrt{q_\parallel^2 - 4m_q^2} + \sqrt{q_\parallel^2}} \right)^2, \quad (3)$$

where again the condition $q_\parallel^2 > 4m_q^2$ is to be satisfied. Given that the Z^0 is a finite width resonance described in vacuum by a relativistic Breit-Wigner function, the contribution from the hadron side to the dimuon production rate in a constant magnetic field can be written, in terms of the above

functions, as

$$H^{\mu\nu} = \sum_{i=\parallel, L} \frac{2 \text{Im } \Pi_i P_i^{\mu\nu}}{(q^2 - M_Z^2 - \text{Re } \Pi_i)^2 + (M_Z \Gamma_Z + \text{Im } \Pi_i)^2}, \quad (4)$$

where Γ_Z is the Z^0 total decay width. Equation (4) corresponds to the sum of the two spectral functions contributing to the process. Notice that since around the Z^0 peak both $\text{Im } \Pi_i$ and $\text{Re } \Pi_i$, $i = \parallel, L$, are proportional to (m_q^2/M_Z^2) , the magnetic field driven modifications are negligible and therefore, the spectral density hardly changes, even for optimistic field strengths of order $(10m_\pi)^2$, possibly achieved at the Large Hadron Collider (LHC), where $m_\pi \sim 0.14$ GeV is the mass of the pion. Therefore, the signature of the presence of a strong magnetic field in the dilepton channel should not be sought in the magnetic-field driven distortion of the Z^0 spectrum.

We now turn to describing the contribution from the lepton side. Recall that, in a magnetic field, charged fermions are not described by plane waves but instead by Ritus functions. The amplitude for dimuon emission is obtained by integration at the lepton vertex in the space-time point y and amounts to computing the overlap integral of the muon-antimuon pair wave functions, which in turn produces energy-momentum conservation factors except in the direction of the confining oscillation around the magnetic field lines. This oscillation can be described using different gauge choices for the vector potential associated to the magnetic field. We work in the Landau gauge 2 (LG2), where the oscillation is one-dimensional along the component y_2 at the lepton vertex. The other two possible gauge choices that describe this confining oscillatory motion transverse to the field are LG1, where the oscillation is also one dimensional along the component y_1 , or the symmetric gauge where the oscillation occurs in the transverse plane (y_1, y_2) . The result is, of course, gauge-choice invariant. Working in LG2, momentum is not conserved in the y_2 direction. In fact, the delta functions that in the absence of a magnetic field would represent momentum conservation in the transverse directions are in practice replaced by a transverse area S in the (y_1, y_2) plane where the process is active. The contribution to the probability of dimuon emission is obtained by squaring the amplitude and averaging over the final-state polarizations. Recall that, when working in the LLL, charged fermions can occupy only one possible polarization state that depends on their charge: positive (negative) charge fermions are polarized along (opposite) to the magnetic field direction. The Ritus wave functions in the LLL already take care of this averaging since the description of the process includes the corresponding spin projection operators. Therefore, the tensor in the lepton side is given by

$$L_{\mu\nu} \sim \tilde{C}_A^2 [p_{\parallel\mu}^- p_{\parallel\nu}^+ + p_{\parallel\nu}^- p_{\parallel\mu}^+ - g_{\mu\nu} [(p^- \cdot p^+)_{\parallel} - m_{\text{muon}}^2]] \times (2\pi)^2 \delta(q_0 - E^- - E^+) \delta(q_3 - p_3^- - p_3^+), \quad (5)$$

where $p_{\parallel\mu}^\pm$ and $E^\pm$ are the parallel components of the muon (−) and antimuon (+) four-momenta and their corresponding energies, respectively, and m_{muon} is the muon mass. In writing Eq. (5) we have anticipated that the largest contribution to the process comes from the axial coupling $\tilde{C}_A$ of the Z^0 with the muon pair.

To obtain the dimuon emission rate per unit time and volume, we recall that, when the process is immersed into a magnetic field, integration over the transverse-momentum coordinates of the phase space is replaced by a sum over Landau levels weighed with a density-of-states factor. In the present case, where only the LLL contributes, this amounts to keeping the factor $(|eB|/2\pi)$ for each of the charged particles in the final state. Contracting the hadron and lepton tensors, and expressing results in terms of the pair invariant mass $M = \sqrt{q_0^2 - q_3^2}$ and rapidity $\eta = (1/2) \ln[(q_0 + q_3)/(q_0 - q_3)]$, we can finally write the expression for the invariant mass distribution of the magnetic field-induced number of dimuons emitted per unit time and volume, $(dN/dV dt)$, as

$$\begin{aligned} & \frac{dN/dV dt}{dM^2 d\phi} \\ &= \tilde{C}_A^2 \left(\frac{|eB|^2}{\pi} \right) \left(\frac{m_{\text{muon}}^2}{M^2} \right) \frac{e^{-a \frac{q_3^2}{|eB|}}}{\sqrt{1 - 4m_{\text{muon}}^2/M^2}} \\ & \times \sum_{i=\parallel, \perp} \frac{|\text{Im } \Pi_i|}{(M^2 - q_\perp^2 - M_Z^2 - \text{Re } \Pi_i)^2 + (M_Z \Gamma_Z + \text{Im } \Pi_i)^2}, \end{aligned} \quad (6)$$

where $a = (1 + 2|e_q/e|)$.

III. DIMUON SPECTRUM PROPERTIES

The invariant-mass distribution, Eq. (6), possesses very distinctive properties. First, it describes the production of muon-antimuon pairs whose direction of motion is highly collimated along lines parallel to the magnetic field, which are in turn perpendicular to the reaction plane. This can be quantified by computing the second harmonic coefficient v_2 of the azimuthal dilepton distribution and recalling that a positive (negative) v_2 corresponds to motion inside (outside) the reaction plane. From Eq. (6), the azimuthal angle ϕ distribution is given by $dN/d\phi \sim e^{-a|\vec{q}|^2/|eB|\cos^2\phi}$, where $|\vec{q}|^2$ is the magnitude of the dilepton-pair momentum which, around the Z^0 peak, can be taken as $|\vec{q}|^2 \sim |eB|$, and where for simplicity we have neglected the dependence on the magnitude squared of the transverse momentum in the denominator of Eq. (6) compared to M_Z^2 . v_2 is thus given by

$$v_2 \sim \frac{1}{2\pi} \int_0^{2\pi} e^{-a\cos^2\phi} \cos 2\phi d\phi = -e^{-\frac{a}{2}} I_1\left(\frac{a}{2}\right) < 0,$$

where I_1 is the Bessel function 1 of the second kind. The negative value of v_2 implies a preferred out-of-plane emission of dileptons. In contrast, the corresponding dilepton emission rate in vacuum is isotropic in the transverse plane. We point out that this anisotropy coefficient has been measured in $\sqrt{s_{NN}} = 5.02$ TeV Pb-Pb collisions by the CMS Collaboration for dileptons originating in the decay of Z^0 and γ^* , who found it compatible with zero, albeit with large uncertainties [51]. Second, Eq. (2) shows that the longitudinal and transverse Z^0 polarizations are proportional to the coefficients C_A^2 and C_V^2 , respectively. Since C_A^2 is an order of magnitude larger than C_V^2 , the longitudinal Z^0 polarization state is preferred

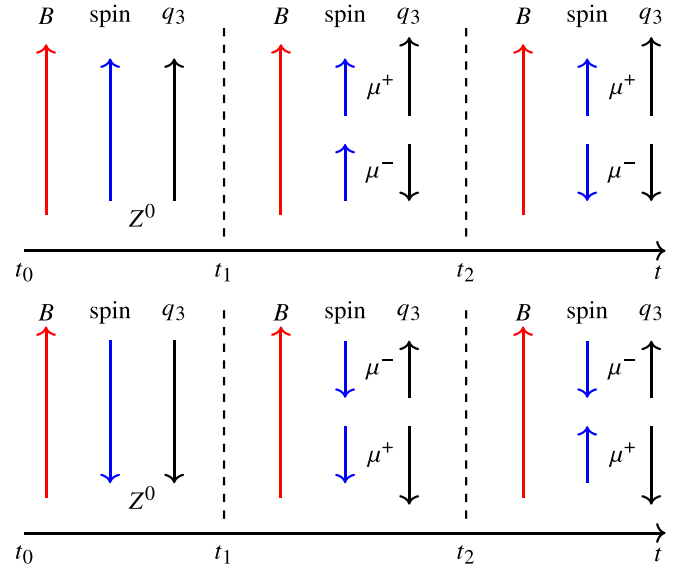


FIG. 3. Schematic representation of the time evolution of the Z^0 decay into a dimuon pair in a strong magnetic field. Z^0 's are produced early after the collision at t_0 and polarized preferably in the longitudinal direction to subsequently decay into a dimuon pair at t_1 by means of the electroweak interaction producing a right-handed antimuon and a left-handed muon. When the Z^0 's move in the direction (opposite direction) of $\vec{B}$, μ^+ 's, being right handed, are produced with larger momentum than μ^- 's; this is depicted in the upper (lower) panel. Overall μ^+ 's are produced with a larger momentum, which should be reflected in their p_T distribution as compared to μ^- . If the field is intense enough, for Z^0 moving in the direction (opposite direction) of $\vec{B}$, the muon (antimuon) makes a spin flip transition to occupy the preferred spin orientation in the LLL at t_2 . This is depicted in the third column of both panels.

over the transverse one. This means that muons and antimuons are produced with their spins pointing in the same direction that the Z^0 is moving. Recall that in this case muons are produced by the electroweak interaction. Given the large mass difference between the Z^0 and the muons, the latter behave as if they were massless, and thus muons are preferably left handed whereas antimuons are preferably right handed. This means that muons move in the opposite direction with respect to that of the antimuons. Given the correlation between muon-antimuon polarization with respect to the direction of motion of the parent Z^0 , altogether these requirements mean that the peak of the antimuon distribution is displaced to a larger p_T value with respect to the peak of the muon distribution, which in turn is displaced to a smaller p_T value. To reconcile this picture with the fact that in the LLL the spin polarization state for negative charge particles is opposite to the field direction, we need to recall that, whereas the electroweak dimuon production can be considered as instantaneous, a short but finite time is required before the spins of the muons and antimuons are influenced by the presence of the magnetic field. This relaxation time is typically of the order of the inverse of the Larmor frequency, namely $\tau \sim m_{\text{muon}}/|eB|$, which shows that as the mass of the produced particle becomes smaller and/or the field strength becomes larger, τ becomes shorter. The time sequence described above is illustrated in Fig. 3. Thus,

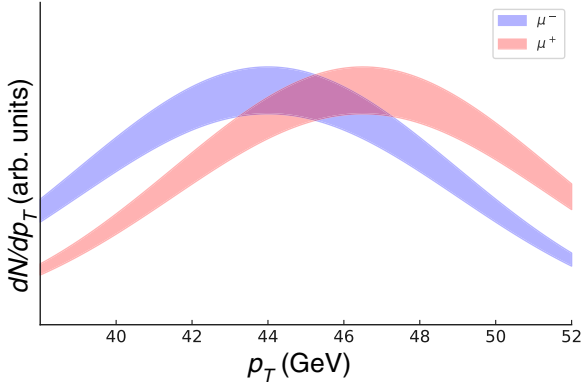


FIG. 4. Schematic representation of the expected p_T distribution of $\mu^\pm$ in high-energy semicentral heavy-ion collisions. The correlation between muon-antimuon polarization with respect to the direction of motion of the parent Z^0 implies that antimuons move with a larger p_T than muons. This should be reflected in an antimuon distribution with a peak shifted to a larger p_T with respect to the muon distribution peak. The bands represent 10% uncertainties associated with the strength of the field, contributions from levels other than the LLL, and the Z^0 p_T distribution from quark distribution functions in the colliding nuclei. In contrast, the muon and antimuon p_T distributions in $p + p$ collisions should show very similar trends.

when $\tau = t_2 - t_1$ is shorter than the time the system spends in the preequilibrium stage, all the above characteristics provide a clear signal for the presence of the magnetic field during dimuon production. The expected transverse momentum distribution for muons and antimuons is qualitatively illustrated in Fig. 4. The bands represent the estimated uncertainty when varying the parameters of the fit to the experimental Z^0 distribution as well as the possible effect of including a few Landau levels on top of the LLL.

IV. DISCUSSION AND CONCLUSIONS

In conclusion, we have shown, for the first time to our knowledge, that a very clear experimental signal of the presence of a strong magnetic field during the earliest stage of a heavy-ion collision can be sought by detecting dimuon pairs whose invariant mass is around the Z^0 peak. The process is anisotropic, and the muon pair is preferably emitted along the field direction. The process also produces an antimuon peak displaced toward larger values of p_T compared to the muon peak. We have also shown that the process does not produce a significant distortion of the Z^0 spectral function. The computation resorts to the LLL approximation. Since muon emission is a sequential process that starts with the description

of a quark-antiquark pair in the presence of a large intensity field, the condition for the validity of LLL dominance is that the quark mass squared is small compared to the field strength. Given that the most abundant quarks in the nucleons making up the colliding nuclei that participate in the initial hard collision are the u and d flavors, the approximation is well justified. Even if the s -quark contribution is included, the approximation is still justified for the kind of field strengths expected at LHC energies. The contribution from levels other than the LLL can also be estimated as follows: recall that in a uniform magnetic field the Ritus eigenstates along the transverse direction contain a Gaussian factor whose argument is the ratio of the relative transverse momentum squared to the field strength $x = q_\perp^2/|eB|$, multiplied by Laguerre polynomials of the same argument. Therefore, the contribution from the n th Landau level scales as $x^n \exp -x$, which for small x is parametrically suppressed by powers of x , justifying the use of the LLL approximation. A more refined analysis certainly calls for the inclusion of at least a few more Landau levels.

The present work provides a general framework to call attention to the experimental search for the signal, drawing from the description of the main features of the process and capturing its main characteristics. It also calls for a more quantitative and refined analysis for a more direct comparison to experimental results. We are currently working in this direction and will report our findings elsewhere.

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DATA AVAILABILITY

No data were created or analyzed in this study.

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