

Quark Coalescence based on a Transport Equation

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We employ the Boltzmann equation for describing hadron production from a quark-gluon plasma (QGP) in ultrarelativistic heavy-ion collisions. We propose resonance formation in quark-antiquark scattering as the dominant meson-production channel, which, in particular, ensures that energy is conserved in the recombination process. This, in turn, facilitates a more controlled extension of hadronization to low transverse momenta (p_T), and to address the experimentally observed transition from a hydrodynamic regime to constituent quark-number scaling (CQNS). Based on input distributions for strange and charm quarks with azimuthal asymmetries, $v_2(p_T)$, characteristic for RHIC energies, we recover CQNS at sufficiently high p_T , while at low p_T a scaling with transverse kinetic energy is found, reminiscent to experiment. The dependence of the transition regime on microscopic QGP properties, i.e. resonance widths and Q -values in the $q + \bar{q} \rightarrow M$ process, is elucidated.

I. INTRODUCTION

Among the surprising experimental results in Au-Au collisions in the first years of operation of the Relativistic Heavy-Ion Collider (RHIC) were several unexpected features of hadron spectra in the intermediate transverse-momentum regime, $p_T \simeq 2 - 6$ GeV [1, 2]. Most notably, a large baryon-to-meson ratio of around one and the so-called constituent quark-number scaling (CQNS) of the elliptic flow coefficient, $v_2(p_T)$, have been found. These features are difficult to explain in hydrodynamic models [3, 4, 5] which provide an excellent description of hadronic observables in the low- p_T region, $p_T \leq 2$ GeV, but lack sufficient yield at intermediate p_T . Quark coalescence models (QCMs) have been used extensively and successfully to describe the baryon-over-meson enhancement and CQNS at intermediate p_T [6, 7, 8, 9, 10]. The basic mechanism underlying QCMs is the recombination of constituent quarks into hadrons at the putative phase boundary (gluonic degrees of freedom are either neglected or converted “by hand” into $q\bar{q}$ pairs). The constituent quarks are identified with the *valence* quarks inside hadrons, and, consequently, the hadron spectra directly reflect upon the partonic spectra and collective dynamics of the Quark-Gluon Plasma (QGP). A major challenge is then to search for (and possibly construct) a unified description of hadron production at low and intermediate p_T . In fact, as demonstrated in a recent PHENIX paper [11], a very general CQNS seems to emerge (in centrality, system size and collision energy) when the scaling is applied to the *kinetic energy* of the produced hadrons.

One of the limitations of conventional QCMs is the collinear and instantaneous approximation of the hadron formation process which limits their applicability to intermediate p_T where the inherent violation of energy conservation is expected to be small. It is also not obvious how to recover an equilibrium limit in this framework which renders the connection to the low p_T (hydrodynamic) regime less transparent. In this work, we propose an alternative realization of the recombination picture by formulating hadron production in terms of a Boltzmann transport equation. Our key dynamical ingredient is hadronic resonance formation via $q\bar{q}$ “annihilation”. Their finite width not only enables to go beyond the collinear limit and explicitly conserve 4-momentum, but also introduces, in principle, a sensitivity to spectral properties of the QGP (e.g., widths and Q values in the fusion process). In the limit of large times, on which we will focus here, we recover an equilibrium formula which reproduces the thermal Boltzmann limit and can be compared to QCM results. We will constrain ourselves to the cases of charm and strange quarks hadronizing within a bulk medium where the input anti-/quark phase distributions are parametrized with collective properties inferred from RHIC data. We will compute both p_T spectra and $v_2(p_T)$ for ϕ , D and J/ψ mesons, and study their scaling properties in connection with the transition from low to intermediate p_T . Our approach bears some similarity with recent work in Ref. [12], which, however, does not address the v_2 problem and utilizes different interactions.

II. BOLTZMANN EQUATION WITH QUARK-ANTIQUARK RESONANCES

Our starting point is the Boltzmann equation for the meson phase-space distribution, f_M ,

$$\left(\frac{\partial f_M}{\partial t} + \vec{v} \cdot \vec{\nabla} f_M\right)(x, p) = -\Gamma/\gamma_p f_M(x, p) + \beta(x, p), \quad (1)$$

where p and x denote 3-momentum and position of the meson M , and Γ its width which in the present work is attributed to 2-body decays into quark and antiquark, $M \rightarrow q + \bar{q}$, and assumed to be constant. The factor

$\gamma_p = E_M(p)/m$ ($m, E_M(p)$: meson mass and energy), is due to Lorentz time dilation (see also Ref. [12]). The surface term on the left-hand-side of Eq. (1) is mostly relevant for high- p_T particles, and vanishes in the infinite volume limit; since our objective here is not a quantitative description of data we neglect it in the following. The relation of the gain term, $\beta(x, p)$, to the underlying microscopic interaction becomes explicit upon integration over phase space,

$$B \equiv \int \frac{d^3 p d^3 x}{(2\pi)^3} \beta(p, x) = \int \frac{d^3 x d^3 p_1 d^3 p_2}{(2\pi)^6} f_q(x, p_1) f_{\bar{q}}(x, p_2) \sigma(s) v_{rel}(p_1, p_2) \quad (2)$$

with $\sigma(s)$ the cross section for the process $q + \bar{q} \rightarrow M$ at center-of-mass (CM) energy squared $s = (p_1^{(4)} + p_2^{(4)})^2$, $p_{1,2}^{(4)}$ the 4-momenta of quark and antiquark, and $f_{q,\bar{q}}$ their phase space distribution functions, normalized as $N_{q\bar{q}} = \int \frac{d^3 x d^3 p}{(2\pi)^3} f_{q,\bar{q}}(x, p)$. Throughout this paper, quarks will be assumed to be zero-width quasi-particles with an effective mass m_q which we treat as a parameter. The intrinsically classical nature of the Boltzmann equation warrants the use of classical distribution functions for all the particles. In addition, we will assume zero chemical potential for all quark species. For the cross section we employ a relativistic Breit-Wigner form,

$$\sigma(s) = g_\sigma \frac{4\pi}{k^2} \frac{(\Gamma m)^2}{(s - m^2)^2 + (\Gamma m)^2}, \quad (3)$$

where $g_\sigma = g_M/(g_q g_{\bar{q}})$ is a statistical weight given in terms of the spin (-color) degeneracy, g_M ($g_{q,\bar{q}}$), of the meson (anti-/quark); k denotes the quark 3-momentum in the CM frame. With $M \rightleftharpoons q + \bar{q}$ being the only channel, it follows that $\Gamma_{in} = \Gamma_{out} = \Gamma$. Detailed balance requires the same Γ in the loss term on the right-hand-side of Eq. (1), thus ensuring a proper equilibrium limit with $\tau = 1/\Gamma$ the pertinent relaxation time. This formulation conserves 4-momentum and applies to all resonances M with masses close to or above the $q\bar{q}$ threshold, i.e., for positive Q value,

$$Q = m - (m_q + m_{\bar{q}}) \gtrsim 0. \quad (4)$$

If the $2 \rightarrow 1$ channel proceeds too far off-shell, i.e., $Q < 0$ and $\Gamma < |Q|$ (e.g., for pions), other processes need to be considered, e.g., $q + \bar{q} \rightarrow M + g$ (which is possible in the present formalism by including the respective cross sections).

Let us now elaborate the equilibrium limit of our approach, by imposing the stationarity condition,

$$0 = \frac{dN_M}{dt}|_{eq} = - \int \frac{d^3 x d^3 p}{(2\pi)^3} \Gamma f_M^{eq}(p) / \gamma_p + B. \quad (5)$$

Introducing the notation $p_M = p_1 + p_2, p_{rel} = p_1 - p_2$ into the gain term B , eq. (2), we find

$$N_M^{eq} = \int \frac{d^3 x d^3 p_M}{(2\pi)^3} f_M^{eq}(p_M, x); \quad f_M^{eq}(p, x) \equiv g(p_M, x) \gamma_p / \Gamma \quad (6)$$

$$g(p_M, x) \equiv \int \frac{d^3 p_{rel}}{8 (2\pi)^3} f_q(x, p_M, p_{rel}) f_{\bar{q}}(x, p_M, p_{rel}) \sigma(s) v_{rel}(p_M, p_{rel}). \quad (7)$$

Eq. (6) represents the large time limit of the Boltzmann equation and is the expression which comes closest to the conventional QCM formula. For hadronization times less or comparable to the relaxation time, τ , the equilibrium limit will not be reached and a short time solution will be appropriate. In that case, the time variable enters explicitly into the final result, reflecting the dynamical nature of the Boltzmann equation.

We have verified numerically that for $\Gamma \rightarrow 0$ Eq. (6) accurately recovers the standard Boltzmann distribution for a meson M at temperature T , if the constraint of a positive Q value is satisfied (for negative Q the $2 \rightarrow 1$ channel is inoperative). This shows that equilibration and energy conservation are closely related in our approach, constituting a significant improvement over previous QCMs.

III. QGP FIREBALL, SPECTRAL PROPERTIES AND TRANSVERSE-MOMENTUM SPECTRA

To compute meson spectra we have to specify the input q and \bar{q} distributions (including their masses and collective properties), as well as the meson-resonance masses and widths. For an exploratory calculation, we will focus on a QGP fireball close to the expected hadronization temperature, $T_c \simeq 170$ MeV (as in previous QCM studies). For simplicity, we will assume a constant (average) homogeneous cylindrical volume with collective transverse expansion velocity with linear radial profile, $\vec{v}_T(r_T) = \beta_0 \frac{\vec{r}_T}{R_T}$ (R_T, β_0 : surface radius and flow velocity), and use the quasi-equilibrium limit, Eq. (6). The quark distributions in the local (thermal) frame take the form

$$f_q(p, x) = \exp[-\gamma_T(E_q - \vec{p} \cdot \vec{v}_T)/T] \quad (8)$$

where $\gamma_T = (1 - v_T^2)^{-1/2}$ and $E_q = (m_q^2 + p^2)^{1/2}$ is the quark energy in the lab frame. In this section we evaluate p_T spectra for central $\sqrt{s_{NN}} = 200$ GeV Au-Au collisions, thus neglecting any azimuthal asymmetries. While the shape of the p_T spectra is affected by β_0 , the total multiplicities must be independent of β_0 . Based on the fireball eigenvolume, $V^* = \int d^3x^*$, corresponding to a system at rest ($\beta_0 = 0$), the volume element in the lab frame follows from the corresponding Lorentz-contraction factor $d^3x_{lab} = d^3x^*/\gamma(x)$. We have verified that this leads to total quark multiplicities independent of β_0 . The eigenvolume is fixed at $V^* \simeq 1500\text{fm}^3$ using a hadron resonance gas in chemical equilibrium at $T \simeq 170$ MeV to reproduce the correct pion rapidity density $dN_\pi/dy(y=0)$ for 0-10% central collisions (including feed-down) [13]. Since thermal rapidity distributions for heavy particles are narrower than for thermal pions, the $dN/dy(y=0)$ values for the former have to be renormalized by a rapidity-width ratio, Γ_y^M/Γ_y^π ($\Gamma_y^\pi=1.8$), to recover the proper equilibrium ratio for $dN_M/dy(y=0)$. In this way, our calculated meson spectra are absolutely normalized.

Based on Eq. (6), the invariant meson p_T spectrum takes the form

$$E_M \frac{dN_M}{d^3p} = \frac{dN_M}{d^2p_T dy} = \frac{E_M}{(2\pi)^3} \int \frac{d^3x^*}{\gamma(x)} g(p, x) \gamma_p / \Gamma \quad (9)$$

with $g(p, x)$ defined in Eq. (7). In the azimuthally symmetric case, 3 integrations can be performed analytically, leaving a 3-dimensional integral to be computed numerically. With a dependence on the transverse radius of the form $(\frac{r_T}{R_T})$, the eigenvolume of the fireball factorizes and enters as an overall factor for p_T spectra. A more complete description of p_T spectra should account for a hard component, which at RHIC energies is expected to become significant at $p_T \gtrsim 2 \div 3$ GeV. Since we here focus on the conceptual aspects of our approach, we will defer the inclusion of a hard component to a future paper.

Let us now specify the resonance parameters. The ϕ meson, which in the vacuum has a mass and width of $m_\phi \simeq 1020$ MeV and $\Gamma_\phi \simeq 4$ MeV, is expected to broaden substantially in hot and dense matter (cf., e.g., Ref. [14]), while mass shifts are more controversial. There are indications from lattice QCD that the ϕ survives in the QGP at moderate temperatures above T_c [15]. We will assume for the ϕ its vacuum mass and a default width of $\Gamma_\phi = 50$ MeV, in connection with a strange quark mass of $m_s = 400$ MeV. The concept of D -meson resonances in the QGP has been implemented in Ref. [16] as a means to understand kinetic charm-quark equilibration at RHIC, with fair success in predicting [17] the most recent data for (decay-) electron suppression and v_2 [18]. Here we will assume D -mesons with a mass and width of $m_D = 1.9$ GeV and $\Gamma_D = 100$ MeV, and c - and u -quark masses of $m_c = 1.5$ GeV and $m_u = 350$ MeV. For the J/ψ we employ $m_{J/\psi} = 3.1$ GeV, $\Gamma_{J/\psi} = 100$ MeV.

Fig. 1 illustrates our results for ϕ (left panel), D and J/ψ (right panel) p_T spectra compared to data in central $\sqrt{s_{NN}} = 200$ GeV Au-Au. The surface flow value of $\beta_0 = 0.55$ has been adjusted to best reproduce the ϕ p_T spectra; it is comparable to the value used in the conventional QCMs [7]. The agreement is fair up to the highest currently measured $p_T \simeq 3$ GeV. The spectra are rather insensitive to variations in Γ and Q (provided that $\Gamma \lesssim Q$), but this will be different for the elliptic flow discussed below. For J/ψ and D spectra we have introduced a charm-quark fugacity, γ_c , which is necessary to match the total c -quark number to the one expected from hard production in primordial nucleon-nucleon collisions [23, 24]. Note that $N_{J/\psi} \propto \gamma_c^2$, with γ_c subject to appreciable uncertainty in the charm cross section as well as due to in-medium effects on open-charm states [25], cf. also Ref. [26].

IV. ELLIPTIC FLOW AND SCALING PROPERTIES

Azimuthal asymmetries in the particle's momentum distributions are defined via Fourier moments of order n ,

$$v_n(p_T) = \langle \cos(n\varphi) \rangle(p_T), \quad (10)$$

where the average is over EdN/d^3p with $d^3p/E = dy dp_T^2 d\varphi$. At midrapidity, the odd moments vanish; in the following we will focus on the elliptic flow, $v_2(p_T)$. As discussed in the introduction, the experimental finding of an approximate CQNS relation for essentially all light hadrons at intermediate p_T ,

$$v_2^h(p_T) = n_Q v_2^q(p_T/n_Q) \quad (11)$$

is naturally reproduced by QCMs and constitutes one of their successes (n_Q : number of constituent quarks in hadron h). As in QCMs, our objective is to study how a given input v_2^q in the quark distributions manifests itself on the hadron level. We adopt a factorized ansatz for the quark v_2^q , as used in previous coalescence studies [7, 8],

$$f_q(p, x, \varphi) = f_q(p, x) \otimes (1 + 2 v_2^q(p_T) \cos(2\varphi)) \quad (12)$$

with $f_q(p, x)$ the distribution function of Eq. (8). Note that this ‘‘local’’ implementation of v_2 (cell by cell in the fireball) neglects space-momentum correlations characteristic for hydrodynamic expansion (more realistic distributions

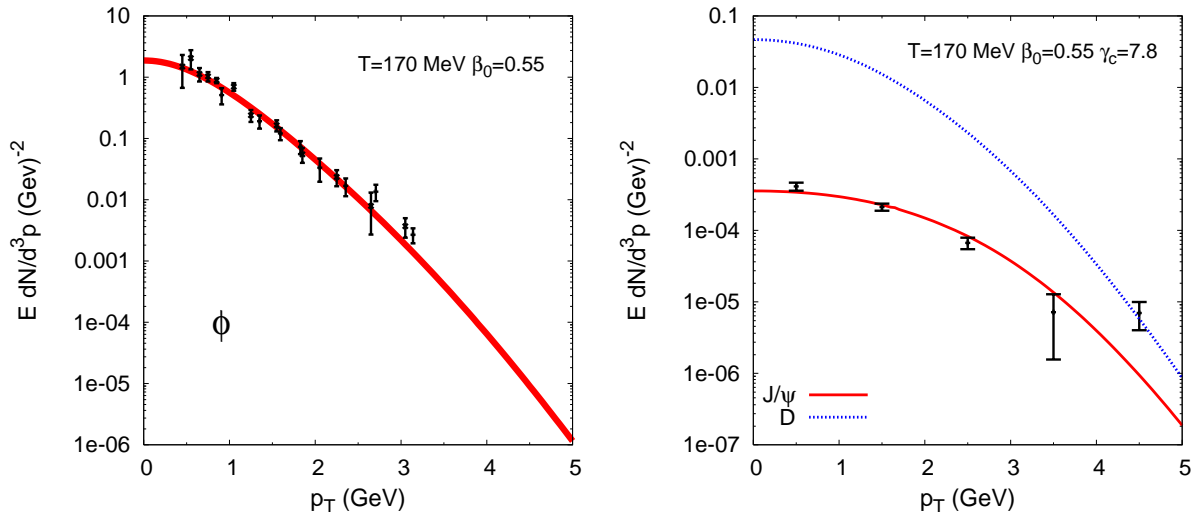


FIG. 1: Left panel: p_T spectra for ϕ resonance formation at $T_c = 170$ MeV within our formalism, compared to data in central Au-Au at RHIC [19, 20, 21]. Right panel: p_T spectra for J/Ψ (lower curve) and D (upper curve) resonance formation, compared to PHENIX J/Ψ data in central Au-Au [22]; $\gamma_c(T = 170 \text{ MeV}) = 7.8$ has been taken from Ref. [25].

will be studied in forthcoming work). For the function $v_2^q(p_T)$ we employ a simplistic ansatz, which, nevertheless, encodes the most important phenomenological features. These are an essentially linear increase at low p_T (as found in hydrodynamic models [3, 4, 5, 27]), and a saturation at intermediate p_T , requiring two parameters: the transition momentum, and the saturation value of v_2 , cf. also Refs. [28, 29]. For light quarks (u, d, s) we choose $p_T^{\text{sat}} = 1.1$ GeV and $v_2^{\text{sat}} = 7\%$, respectively, which gives a reasonable schematic representation of the experimentally found scaling function [11]. Similar features are borne out of parton cascade models [30, 31]. For charm quarks, we take guidance from Langevin simulations for heavy quarks in a thermal background [17, 32]. Also here, a plateau value of $v_2^{\text{sat}} = 7\%$ is within the uncertainty of current data, while the transition momentum is shifted to a higher value, $p_T^{\text{sat}} \simeq 2$ GeV, which is compatible with the mass ordering effect in hydrodynamic calculations [27]. We have verified that our schematic parametrization is consistent with a *parton-level* KE_T scaling for s and c quarks (characteristic for a hydrodynamically expanding parton phase), where $KE_T = m_T - m$ denotes the transverse kinetic energy of a particle ($m_T = \sqrt{m^2 + p_T^2}$). v_2 saturation, implying deviation from hydrodynamics, signals, of course, incomplete thermalization and the transition to a kinetic regime. In Fig. 2 we summarize our results for the $v_2(p_T)$ of ϕ and J/Ψ mesons. The dependence of v_2^M on Γ is very weak, changing by less than 5% when varying the ϕ width over the range $\Gamma = 20 - 400$ MeV. The dependence of v_2^M on the Q value for the resonances is more pronounced as indicated by the different curves which have been obtained by varying the underlying quark masses (e.g., for the ϕ , $Q = 220$ MeV implies $m_s = (m_\phi - Q)/2 = 400$ MeV). For any value of Q , CQNS, Eq. (11), is recovered in our approach at sufficiently high p_T where the production mechanism for an on-shell meson requires two constituent quarks with essentially collinear momenta, $p_q \sim p_{\bar{q}} \sim p_T/2$. Consequently, the meson will inherit the full azimuthal asymmetry imparted by its constituents. As is well known, for collinear production with $v_2^q \ll 1$, using Eq. (12), the product $f_q f_{\bar{q}}$ recovers the scaling relation, Eq. (11).

Deviations from collinearity are expected to induce correction terms involving Q/p_T . Since 4-momentum is conserved in our approach these corrections can be quantified. In particular at low p_T there is no kinematical constraint enforcing collinearity of the reaction. For $Q \rightarrow 0$, we accurately reproduce CQNS at all p_T , with the meson v_2 exhibiting twice the asymptotic value of v_2^q reached at exactly twice the quark transition momentum. For increasing Q , however, the convergence to the limiting value is delayed to higher p_T , together with a reduction of the v_2 at low p_T (resembling the effect of a larger mass).

We finally return to the question of a universal CQNS. A recent analysis of the PHENIX collaboration [11] confirmed that CQNS for hadron p_T spectra is only satisfied at sufficiently high p_T , but the data exhibit a remarkably universal CQNS scaling when plotted vs. kinetic energy of the hadrons, $KE_T = m_T - m$. At low p_T , KE_T scaling is suggestive for hydrodynamic behavior [33], followed by a transition to a kinetic regime where the v_2 levels off. In Fig. 3 we summarize the CQNS properties as computed in our approach. The left panel displays the v_2 for ϕ , D , and J/Ψ when scaled in p_T . The 3 curves are different at low p_T , converging at momenta as high as $p_T \gtrsim 5$ GeV. However,

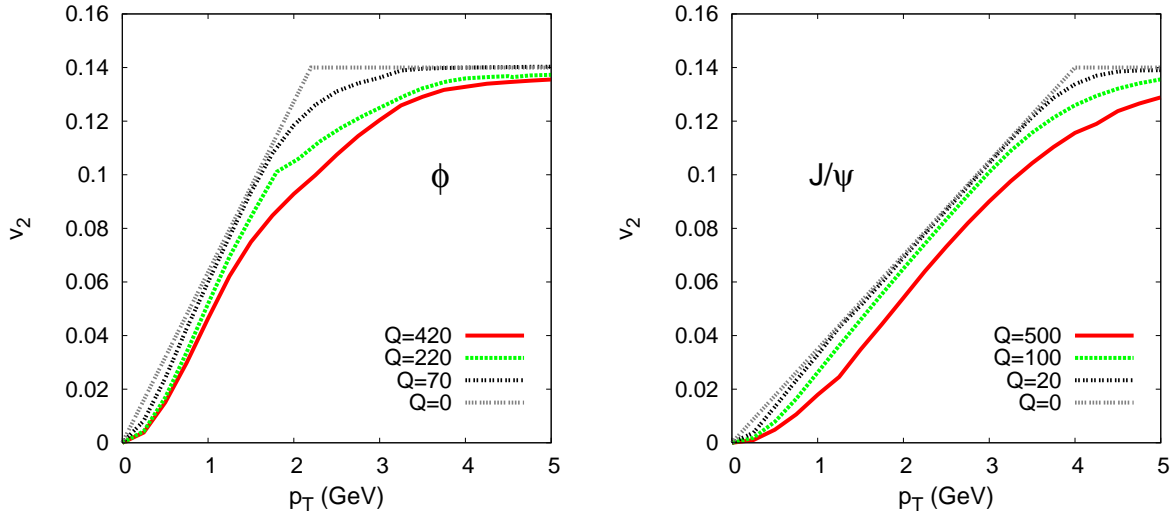


FIG. 2: Elliptic flow for ϕ (left panel) and J/Ψ (right panel) for 4 Q values, $Q_\phi = 420, 220, 70, \sim 0$ MeV and $Q_{J/\Psi} = 500, 100, 20, \sim 0$ MeV, from bottom up. The $Q \rightarrow 0$ limits (upper curves) accurately reproduce the scaled quark v_2 .

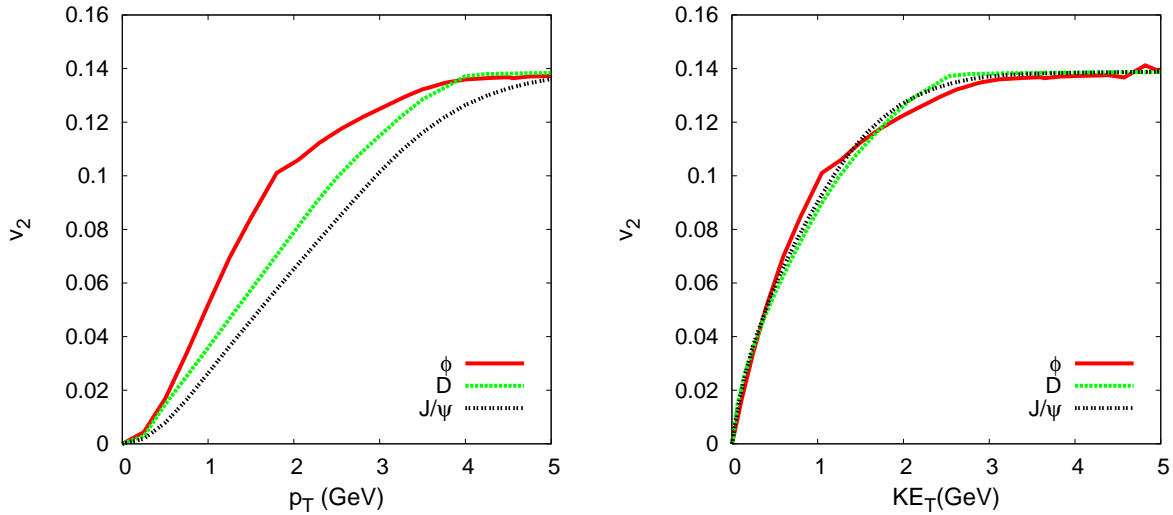


FIG. 3: Left panel: v_2 for ϕ , D and J/Ψ (top to bottom) using the default values for Q and Γ ; universal CQNS emerges only at high p_T . Right panel: v_2 for ϕ , D and J/Ψ vs. kinetic energy, KE_T , of the meson; a universal scaling appears to emerge.

when the scaling is applied in KE_T (right panel of Fig. 3), the curves essentially coincide over the entire energy range, in qualitative agreement with experiment. While suggestive for the transition from the hydrodynamic to the kinetic regime, it remains to be scrutinized in how far this result depends on the uncertainties in the underlying parametrization of the quark v_2 's. As discussed above, the latter have been approximated by hydro-like behavior at low p_T with a sharp transition to an asymptotic value of 7% at a momentum as estimated from transport calculations and data. An interesting point is that a KE_T scaling implemented at the partonic level appears to manifest itself as KE_T scaling at the meson level *if* the Q value in the $q + \bar{q} \rightarrow M$ reaction is not too large (typically below 300 MeV, recall Fig. 2). Another critical assumption is equality of the saturation value for v_2^q for $q = u, s, c$, which, in turn, is suggestive for full c -quark collectivity in a strongly coupled QGP (sQGP). We finally note that, especially for the J/ψ , the contribution from regeneration processes (as calculated here) is not expected to account for the yield at $p_T \gtrsim 4$ GeV (cf. right panel of Fig. 1), which implies deviations from CQNS at high momenta.

V. CONCLUSIONS

We have proposed a reformulation of quark coalescence approaches based on a transport equation by implementing resonant quark-antiquark interactions to form mesons as the key microscopic mechanism. Our approach improves previous coalescence models by including explicit energy conservation in the meson production process and a well-defined thermal equilibrium limit. In the present exploratory study we have focused on a quasi-stationary scenario by considering the meson formation process in a mixed phase at constant temperature, $T \simeq T_c$. We recover the empirically found constituent quark-number scaling (CQNS) in the meson v_2 at high momenta, while the improvements allow to extend the description to the low-momentum regime. We have quantified deviations from CQNS in p_T in terms of underlying spectral properties of the meson resonances, i.e., their widths and Q values. Larger values of the latter are found to shift the onset of CQNS to higher momenta. An interesting result of our study is that, for input (c - and s -) quark distributions satisfying kinetic-energy scaling, a KE_T scaling is recovered at the meson level (ϕ , D , J/ψ) provided the Q -values are positive and not too large. A universal KE_T scaling of hadron v_2 has recently been established experimentally [11]. To consolidate our results, the sensitivity to the input $v_2(p_T)$ at the quark level (including space-momentum correlations) needs to be scrutinized [10]. Future work should also address the problem of baryon formation (e.g., within a quark-diquark picture [12]), an explicit time evolution (before and after T_c) as well as hadron channels with negative Q values (e.g., for Goldstone bosons π and K). While our framework does not address the hadronization problem of gluons, we believe that it could lead to useful insights into systematic features of hadron production over a rather wide range of momenta.

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