Production of $\phi$ and $\Omega^-$ at RHIC in the Recombination Model

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The recombination model is applied to the production of $\phi$ and $\Omega^-$ at all $p_T$ in central Au+Au collisions. Since no light quarks are involved in the hadronization, those hidden-strange particles present a clean slate for the study of the role of strange quarks in large-$p_T$ physics. We find that shower $s$ quarks have negligible effect for $p_T < 6$ GeV/$c$, in which range the thermal $s$ quarks make the dominant contributions to the formation of $\phi$ and $\Omega^-$. We show that the same effective temperature of the $s$ quarks is responsible for the shape of the spectra of both $\phi$ and $\Omega^-$. We predict that the ratio of $\Omega^-$ to $\phi$ will show a peak at $p_T \approx 6$ GeV/$c$ due to the effect of the hard partons. We also give reasons on the basis of the $p_T$ dependence that $\phi$ cannot be formed by means of $K^+K^-$ coalescence.

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I. INTRODUCTION

The production of strange particles has always been a subject of great interest in heavy-ion collisions because of their relevance to possible signatures of deconfinement and flavor equilibration [1, 2]. Strangeness enhancement that has been observed at various colliding energies is a phenomenon associated with soft particles in the bulk matter [3, 4]. At high transverse momentum ($p_T$), on the other hand, the production of jets does not favor strange particles, whose fragmentation functions are suppressed compared to those for non-strange particles. At intermediate $p_T$ range between the two extremes the $p_T$ distribution depends sensitively on both the strangeness content and the production mechanism. It has been shown that in that $p_T$ range the spectra of $\pi, K$, and $p$ can be well described by parton recombination [5]. In this paper we study the production of $\phi$ and $\Omega^-$, both of which consist of only strange quarks. Without the participation of the non-strange quarks, they present a clean platform for the examination of the transverse momentum spectra of the $s$ and $\bar{s}$ quarks. Our aim is to learn about the transition from the enhanced thermal quarks to the suppressed shower quarks in the strange sector.

Since the hidden-strange particles ($s\bar{s}, ss$) are expected to have small hadronic cross sections due to the OZI rule [6, 7, 8], $\phi$ and $\Omega^-$ are less likely to be affected by final-state interaction with co-movers in heavy-ion collisions, compared to kaons and hyperons. Their spectra should therefore reveal more directly their formation mechanism. We shall apply the recombination model, as in [9], and predict the shapes of their spectra beyond the intermediate $p_T$ range where data do not yet exist. We can also calculate the $\Omega/\phi$ ratio and show how different it is from the $p/\pi$ ratio that has been the definitive signature of parton recombination [9, 10, 11]. If the prediction of the peaking of the $\Omega/\phi$ ratio at $p_T \approx 6$ GeV/$c$ is verified by experiments, it will be another piece of evidence in support of recombination at high $p_T$.

II. FORMULATION OF THE PROBLEM

We shall assume that all hadrons produced in heavy-ion collisions are formed by recombination of quarks and/or antiquarks, the original formulation of which is given in [12, 13]. Later improvements are described in [15] for hadronic collisions, in [16] for $pA$ collisions and in [3] for AA collisions. For any colliding system the invariant inclusive distribution of a produced meson with momentum $p$ in a 1D description of the recombination process is

$$p^0 dN_M / dp = \int dp_1 dp_2 F_{qq'}(p_1, p_2) R_M(p_1, p_2, p), \quad (1)$$

and for a produced baryon

$$p^0 dN_B / dp = \int dp_1 dp_2 dp_3 F_{qq'q''}(p_1, p_2, p_3) R_B(p_1, p_2, p_3, p). \quad (2)$$

The properties of the medium created by the collisions are imbedded in the joint quark distributions $F_{qq'}$ and $F_{qq'q''}$. The recombination functions (RF) $R_M$ and $R_B$ depend on the hadron structure of the particle produced, but not on the medium out of which quarks hadronize. In
the thermal source as (1) and (2), we obtain the simple algebraic expressions for $S$ expressed as.

where $y_i = p_i/p$, and $g_M$ and $g_A$ are statistical factors. $G_M$ and $G_B$ are the non-invariant probability densities of finding the valons with momentum fractions $y_i$ in a meson and a baryon, respectively.

Equations (1) and (2) apply for the produced hadron having any momentum $p$. We consider $p_\perp$ only in the transverse plane, and write $p_T$ as $p$ so that $dN/dp_T dp_T$ becomes $(p_T^0)^{-1}$ times the right-hand sides of Eqs. (1) and (2). For $φ$ and $Ω^-$, production, $F_{qq}$ is $F_{ss}$, and $F_{qq'q''}$ is $F_{ss}$. The RFs are very narrow in momentum space, since both $φ$ and $Ω^-$ are loosely bound systems of the constituent $s$ quark. We shall approximate $G_φ$ and $G_{Ω^1}$ by $δ$-functions:

$$G_φ(y_1, y_2) = δ(y_1 - 1/2) δ(y_2 - 1/2), \quad G_{Ω^1}(y_1, y_2, y_3) = δ(y_1 - 1/3) δ(y_2 - 1/3) δ(y_3 - 1/3)$$

and set $g_A = 3/4$ and $g_M = 1/2$ for $J = 1$ and $J = 3/2$, respectively. Using these in Eqs. (3) and (4), and then in (1) and (2), we obtain the simple algebraic expressions

$$\frac{dN_φ}{dp} = 3 \frac{F_{ss}}{4p_T^0} F_{ss}(p/2, p/2), \quad \frac{dN_A}{dp} = \frac{1}{2p_T^0} F_{ss}(p/3, p/3, p/3),$$

where $p_0 = (m^2 + p_T^2)^{1/2}$, $m$ being the mass of $φ$ or $Ω^-$, as the case may be.

In (1), we have described how the joint parton distributions can receive contributions from the thermal ($T$) and shower ($S$) sources. In a schematic way they can be expressed as

$$F_{ss} = T T T + T S + S S, \quad F_{ss} = \kappa (T T T + T T S + T S S + S S S),$$

where showers from more than 1 jet are neglected. In Eq. (10) the multiplicative factor $κ$ is added to the $ss$ distribution to allow for the possible constraint arising from the competition among various channels of hadronization that can limit the number of $s$ quarks available for the formation of $Ω^-$. We parameterize the invariant $s$ quark distribution in the thermal source as

$$T(p_1) = C_s p_1 \exp(-p_1/T_s),$$

where $C_s$ and $T_s$ are two parameters to be determined by fitting the low-$p_T$ data; they are expected to be different from those in the non-strange sector. The distribution of shower $s$ quark in central Au+Au collisions is, as in (9),

$$S(p_2) = ξ \sum_i \int_{k_0}^{∞} dk k f_i(k) S_i^s(p_2/k),$$

where $S_i^s$ is the shower parton distribution (SPD) for an $s$ quark in a shower initiated by a hard parton $i$, $f_i(k)$ is the transverse-momentum ($k$) distribution of hard parton $i$ at midrapidity in central Au+Au collisions, and $ξ$ is the average fraction of hard partons that can emerge from the dense medium to hadronize. As in (10), $f_i(k)$ is taken from Ref. [15], $k_0$ is set at 3 GeV/c, and $ξ$ is found to be 0.07.

Equations (11)-(14) completely specify the problem. There are 3 parameters to vary to fit the $p_T$ spectra of $φ$ and $Ω^-$: they are $C_s$, $T_s$ and $κ$. All aspects of the semi-hard shower parts have been fixed by previous studies [16]. The three parameters are all related to the soft thermal parts, the properties of which, we have stated from the outset, are to be determined phenomenologically.

III. RESULTS

We first vary $C_s$ and $T_s$ to fit the data on the $p_T$ distribution of $φ$ in central Au+Au collisions at $\sqrt{s} = 200$ GeV [17]. For $p_T$ up to 3 GeV/c, which is the extent to which data exist, the entire distribution can be accounted for by the recombination of the thermal partons only. The thermal-thermal component is shown by the dashed line in Fig. 1. The values of $C_s$ and $T_s$ used for the fit are

$$C_s = 8.64 (\text{GeV/c})^{-1}, \quad T_s = 0.385 \text{ GeV/c.}$$

The contributions from thermal-shower (dash-dot line) and shower-shower (line with crosses) recombination do not become important until $p_T ≈ 7$ GeV/c. The solid line in Fig. 1 indicates the sum of all contributions. The dominance of the thermal component for $p_T < 6$ GeV/c is due to the fact that the production of $s$ quarks in a shower is highly suppressed [12]. For that reason the role of the shower partons in the formation of $φ$ in the intermediate $p_T$ region is insignificant compared to that for the production of pions [16]. Shower-shower recombination that can be related to fragmentation can become important, but not until $p_T > 8$ GeV/c.

Next, for the production of $Ω^-$ there is only one parameter $κ$ to vary. Again, the recombination of thermal partons only dominates throughout the whole $p_T$ region shown in Fig. 2. All other contributions that involve the participation of at least one shower parton are increasingly negligible with increasing number of shower $s$ quark. The recombination of $SSS$ can be identified with the fragmentation of hard partons to $Ω^-$ in the same sense that the recombination of only the shower $u d d$ in a jet gives the proton by fragmentation [16], even though the fragmentation function for $Ω^-$ does not exist. Evidently, the contribution to the $Ω^-$ spectrum from fragmentation is negligible until $p_T$ is extremely large.

The value of $κ$ used in the fit of the data [17] is

$$κ = 0.037.$$

(14)

It affects only the overall normalization of the spectrum, not the relative magnitudes of the different components.
FIG. 1: Transverse momentum distribution of $\phi$ in central Au+Au collisions. Data are from [19]. The solid line is the sum of the three contributions: $TT$ (dashed line); $TS$ (dash-dot line); $SS$ (line with crosses).

FIG. 2: Transverse momentum distribution of $\Omega^{-}$ in central Au+Au collisions. Data are from [20]. The heavy solid line is the sum of the four contributions: $TTT$ (dashed line); $TTS$ (dash-dot line); $TSS$ (light solid line); $SSS$ (line with crosses).

Thus our predictive power is in the shape of $\Omega^{-}$ spectrum, not in the normalization. The agreement with the data in Fig. 2 is clearly excellent. Apart from the different powers of $p_1$ in Eq. (11) that appear in [11] and [10], the shape is controlled by $T_s$, which is common for both $\phi$ and $\Omega^{-}$ production. The low level of $\Omega^{-}$ production is a phenomenological fact that is embodied in the smallness of $\kappa$ in Eq. (14). The underlying physics is most likely the demand for $s$ quarks in the formation of lower-mass particles, such as kaons, hyperons and $\phi$ meson, so that their availability for forming the higher-mass $\Omega^{-}$ is significantly reduced. To consider all channels of hadronization simultaneously would require a study similar to those carried out in [21, 22, 23], which is beyond the scope of this work, where our emphasis is on the $p_T$ dependences of $\phi$ and $\Omega^{-}$.

Having obtained $\phi$ and $\Omega^{-}$ spectra, let us present their ratio $R_{\Omega/\phi}(p_T)$, shown in Fig. 3. The dashed line is the ratio of only the thermal contributions. It is linearly rising because, on account of Eqs. (9) and (10), it is proportional to $p_1$ at large $p_T$ after the common exponential terms of the two spectra are cancelled. The solid line shows the effect when the thermal-shower recombination is taken into account. Although the effect of shower partons on $\Omega^{-}$ is small, the effect on $\phi$ above $p_T = 5$ GeV/c is substantial enough to turn the dashed line in Fig. 3 into a peak at around 6 GeV/c. Confirmation of this peak by future experimental data would be a definitive validation of the theoretical approach taken here and, in particular, the importance of thermal-shower recombination.

IV. DISCUSSION

The strange sector differs from the non-strange sector in several ways. Firstly, the values of $C_s$ and $T_s$ given in Eq. (13) are different from

$$C = 23.2 \text{ (GeV/c)}^{-1}, \quad T = 0.317 \text{ GeV/c}, \quad (15)$$

for the non-strange quarks $s$. Strangeness enhancement refers to the excess of $s$ quarks in $AA$ collisions compared to the scaled $pp$ collisions, but the total number of $s$ quarks is still much less than that of light quarks in nuclear collisions. That difference is reflected in $C_s$ being much less than $C$. The inverse slope $T_s$ is, however, higher than $T$, since hydrodynamical flow gives the more massive $s$ quark a larger mean transverse momentum than it does the light quarks. The combined effect
of larger $T_s$ and weaker shower $s$ quarks results in the thermal partons becoming a dominant contributor to the formation of $\phi$ and $\Omega^-$ over a much wider range of $p_T$ than in the case of non-strange quarks. For pion production the thermal-shower recombination becomes important for $p_T > 3$ GeV/c, whereas for $\phi$ it is not until $p_T > 7$ GeV/c.

The above comment cannot be checked directly by experiments. However, the predicted ratio of $\Omega^-$ to $\phi$ can be tested. The rising portion of that ratio in Fig. 3 is much broader than that for the $p/\pi$ ratio [24], and is an indication of the dominance of thermal-thermal recombination. The overall shape of $R_{\Omega}/\phi$ is very different from that of $R_{p/\pi}$. The latter peaks at $p_T \approx 3$ GeV/c with an experimental value exceeding 1, whereas the former is predicted to peak at $p_T \approx 6$ GeV/c with a value less than 0.4. Neither $R_{p/\pi}$ nor $R_{\Omega}/\phi$ exhibits properties that can be associated with fragmentation. The difference between them reflects the differences between $u, d$ and $s$ quarks on the one hand, and between non-strange and hidden-strange hadrons on the other.

Within the framework of recombination it is possible for us to examine the reality of $\phi$ formation through $K^+ K^-$ coalescence, which is a mechanism that has been advocated in certain models [25, 26]. If $H_{K}(p_T)$ denotes the invariant inclusive $p_T$ distribution, $p^0 dN_h/dp_T$, of hadron $h$ at $y = 0$ in heavy-ion collisions, then the coalescence process of $K^+ + K^- \rightarrow \phi$ implies by use of Eq. (1) that

$$H_{\phi}^{[KK]}(p_T) \propto H_{K}^{2}(p_T/2)$$

(16)

apart from a multiplicative constant associated with the RF. On the other hand, if $\phi$ is produced by $s\bar{s}$ recombination as we have done here, then the same procedure yields

$$H_{\phi}^{[s\bar{s}]}(p_T) \propto T_s^{2}(p_T/2),$$

(17)

where the dominance of the thermal parton recombination is used. Thus it is a matter of comparing $H_{K}(p_T)$ with $T_s(p_T)$, which are the invariant distributions of the entities that recombine. Since a kaon is formed by $s\bar{s}$ recombination, where $q$ denotes either $u$ or $d$, whose thermal distribution $T(p)$ is characterized by $C$ and $T$ shown in Eq. (15), the exponential part of $H_{K}(p_T)$ must have an inverse slope $T'$ that is between $T$ and $T_s$, i.e., $e^{-p/T'}$ with $2/T' = T^{-1} + T_s^{-1}$. That is to be compared to the exponential part of $T_s^{2}(p_T/2) = T_s(p_T)$, which is $e^{-p/T_s}$ [cf. Eq. (11)]. In view of our good fit of the $\phi$ data in Fig. 1, we conclude that an alternative fit using $H_{K}(p_T)$ characterized by $T'$ would fail. It therefore follows from the consideration of the $p_T$ dependence alone that $\phi$ cannot be formed by $K^+ K^-$ coalescence. This conclusion is consistent with that of [11] based on the centrality independence of the $\phi/K^-$ ratio.

In summary, we have shown that both $\phi$ and $\Omega^-$ are formed mainly by the recombination of thermal $s$ quarks, have reproduced the shape of the $p_T$ distribution of $\Omega^-$ from a fit of that of $\phi$, and have made a prediction of the ratio of $\Omega^-$ to $\phi$ that peaks at $p_T \approx 6$ GeV/c. Thermal-shower recombination does not become important until $p_T$ is at around that peak, and parton fragmentation does not dominate until $p_T$ is much higher. Finally, we have shown how the $p_T$ dependence disfavors the formation of $\phi$ by $K^+ K^-$ coalescence.

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