Hadronization in Nuclear Environment

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Abstract: We present a space-time description of hadronization of highly virtual quarks originating from a deep-inelastic electron scattering (DIS). Important ingredients of our approach are the time- and energy–dependence of the density of energy loss for gluon radiation, the Sudakov's suppression of no radiation, and the effect of color transparency, which suppresses final state interaction of the produced colorless wave packet. The model is in a good agreement with available data on leading hadron production off nucleons and nuclei. The optimal energy range for study of the hadronization dynamics with nuclear target is found to be a few tens of GeV, particularly energies available in the experiment HERMES.

A quark originated from a hard process, converts into colorless hadrons owing to confinement. Lorentz dilation stretches considerably the duration of this process. While hadrons carry a little information about the early stage of hadronization, a nuclear target, as a set of scattering centres, allows us to look inside the process at very short times after it starts. The quark-gluon system produced in a hard collision, interacts while passing through the nucleus. This can yield precious information about the structure of this system and the space-time pattern of hadronization.

The modification of the quark fragmentation function in nuclear matter was considered for high-p_T hadron production in [1, 2], for deep-inelastic lepton scattering in [3, 4, 5], and for hadroproduction of leading particles on nuclei in [6, 7]. The data are usually presented in the form,

$$R_{A/N} = \frac{D_{eff}(z_h, p_T)}{A\ D(z_h, p_T)}, \tag{1}$$

where $D(x, p_T)$ and $D_{eff}(x, p_T)$ are the quark fragmentation function in vacuum and in a nucleus, respectively.

We treat hadronization of a highly virtual quark, perturbatively a gluon bremsstrahlung and the deceleration of the quark as a result of radiative energy loss. We assume that subsequent hadronization of the radiated gluons, which includes the nonperturbative stage, does not affect the energy loss of the quark.

The radiation of a gluon takes the time

$$t_r \approx \frac{2\nu}{k_T^2}\alpha(1-\alpha)\ . \tag{2}$$

This expression follows from the form of the energy denominator corresponding to a fluctuation of a quark of energy ν into a quark and a gluon, having transverse momenta k_T and relative shares of the initial light-cone momentum α and $1-\alpha$, respectively. If one calculates radiated energy taking into account condition (2) one arrives at the density of energy loss per unit of length, which turns out to be energy and time independent [8] like in the string model.

In the case of inclusive production of leading particles at $z_h \to 1$, however, energy conservation forbids the radiation of gluons with energy greater than $(1-z_h)\nu$. Then, the time dependence of the radiative energy loss can be written as

$$\Delta E_{rad}(t) = \int_{\lambda^2}^{Q^2} dk_T^2 \int_0^1 d\alpha\ \alpha\nu \frac{dn_g}{d\alpha dk_T^2} \Theta(1-z_h-\alpha)\ \Theta(t-t_r)\,, \tag{3}$$

where $dn_g/d\alpha dk_T^2 = \epsilon/\alpha k_T^2$ represents the distribution of the number of gluons. The factor $\epsilon = 4\alpha_S(k_T^2)/3\pi$.

Although soft hadronization is usually described in terms of the string model, we model it by radiation as well, choosing the bottom limit λ^2 in (3) small. We fix the QCD running coupling $\alpha_S(k_T^2) = \alpha_S(k_0^2)$ at $k_T^2 \leq k_0^2$, in the region which is supposed to be dominated by nonperturbative effects. the parameter $k_0 \approx 0.7\ GeV$ is chosen to reproduce the density of energy loss for radiation of soft gluons ($k_T^2 \leq k_0^2$) corresponding to the string tension, $dE/dt = \rangle \approx 1\ GeV/fm$. This value of k_0 is consistent with the transverse size of a string corresponding to the gluon-gluon correlation radius, $R_c \sim 1/k_0 \sim 0.3\,\mathrm{fm}$ suggested by QCD lattice results.

After integrating eq. (3) in the soft radiation approximation, we get

$$\Delta E_{rad}(t) = \frac{\epsilon}{2} t(Q^2-\lambda^2)\Theta(t_1-t) + \epsilon\nu(1-z_h)\Theta(t-t_1) +$$
$$\epsilon\nu(1-z_h)\ln\left(\frac{t}{t_1}\right)\Theta(t-t_1)\Theta(t_2-t_1) + \epsilon\nu(1-z_h)\ln\left(\frac{Q^2}{\lambda^2}\right)\Theta(t-t_2) \tag{4}$$

Here we have set $t_1 = (1-z_h)/x_{Bj}m_N$ and $t_2 = t_1\ Q^2/\lambda^2$, where x_{Bj} is the Bjorken variable.

Eq. (4) shows that for $t \leq t_1$, the density of energy loss is constant, $dE/dt = -\epsilon Q^2/2$, exactly as in the case with no restriction on the radiated energy [4, 5]. At longer time intervals, $t > t_1$ more energetic gluons can be radiated and the restriction $\alpha < 1 - z_h$ becomes important. As a result, the density of energy loss slows down to $dE/dt = -\epsilon\nu(1-z_h)/t$ which is a new result compared to what was known in the string model. At still longer $t > t_2$, no radiation is permitted, but obviously a color charge cannot propagate a long time without radiating which must be suppressed by a Sudakov's type formfactor. Assuming a Poisson distribution for the number of emitted gluons we get the formfactor, $F(t) = \exp\left[-\tilde{n}_g(t)\right]$, where $\tilde{n}_g(t)$ is the number of non radiated gluons,

$$\tilde{n}_g(t) = \epsilon\left[\frac{t}{t_1} - 1 - \ln\left(\frac{t}{t_1}\right)\right]\Theta(t-t_1)\,. \tag{5}$$

In order to calculate a time interval for the leading hadron production (or, better, a colorless ejectile which does not loose energy anymore), one needs a model of hadronization and of color neutralization. In the large N_c limit, each radiated gluon can be replaced by a $q\bar{q}$ pair, and the whole system can be treated as a system of color dipoles. It is natural to assume that the leading (fastest) hadron originates from a $q\bar{q}$ dipole made of the leading quark and of the antiquark coming from the last emitted gluon. This dipole is to be projected into the hadron wave function, $\Psi(\beta, l_T)$, where β and $1-\beta$ are the relative shares of the light-cone momentum carried by the quarks, and l_T is the relative transverse momentum of the quarks. The result of this projection leads to the fragmentation function of the quark into the hadron, which reads

$$D(z_h) = \int_0^\infty dt W(t, \nu, z_h) \ , \tag{6}$$

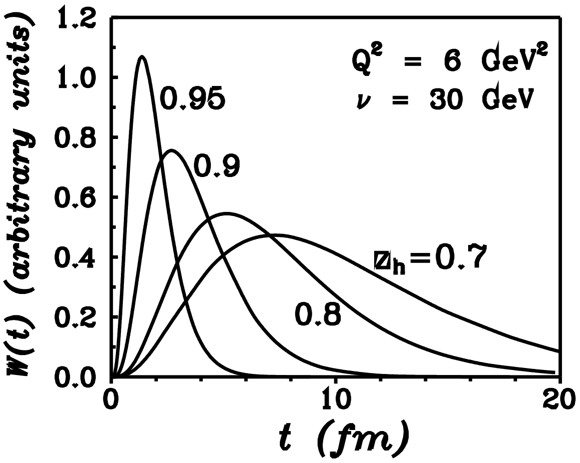
where $W(t, \nu, z_h)$ is a distribution function of the leading hadrons over the production time t.

$$\begin{aligned} W(t,\nu,z_h) \ \ \propto \ \ & \int_0^1 \frac{d\alpha}{\alpha}\, \delta\left[\alpha - 2\left(1 - \frac{z_h \nu}{E_q(t)}\right)\right] \int \frac{dk_T^2}{k_T^2}\, \delta\left[k_T^2 - \frac{2\nu}{t}\alpha(1-\alpha)\right] \times \\ & \int dl_t^2\, \delta\left[l_T^2 - \frac{9}{16}k_T^2\right] \int_0^1 d\beta \delta\left[\beta - \frac{\alpha}{2-\alpha}\right] |\Psi_h(\beta, l_T)|^2 . \end{aligned} \tag{7}$$

Here the quark energy $E_q(t) = \nu - \Delta E_{rad}(t)$. We have chosen a hadronic wave function in the light-cone representation which satisfies the Regge end-point behaviour, $|\Psi_h(l_T^2, \beta)|^2 \propto \sqrt{\beta}\sqrt{1-\beta}(1 + l_T^2 r_h^2/6)^{-1}$, where r_h is the charge radius of the hadron.

Fig.1 shows function $W(t)$ for several values of z_h and exhibits the approximate $(1 - z_h)\nu$-scaling of the mean production time, $t_{pr} = \int dt\ t\ W(t)$, which depends weakly on Q^2 and vanishes at $z_h \approx 1$.

Our predictions for the fragmentation function $D(z_h)$ depicted in Fig. 2 nicely agrees with the EMC data [9].



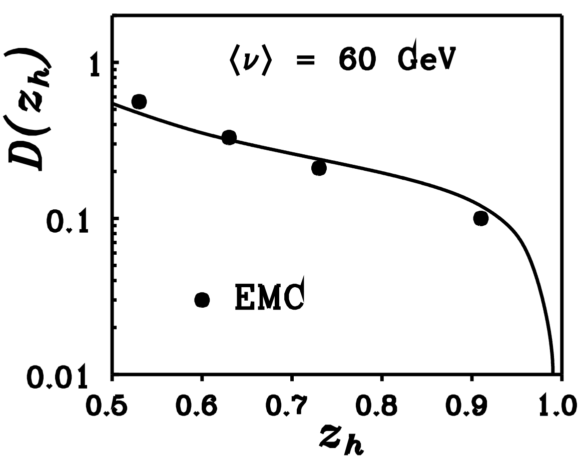
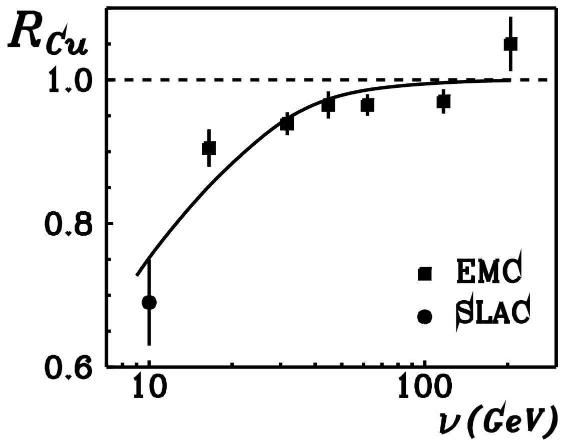


Figure 1: *Distribution of the hadron production time at* $\nu = 30$ *GeV*, $Q^2 = 6$ *GeV*^{2} *and* $z_h = 0.70 \div 0.95$

Figure 2: *Comparison of our prediction for* $D(z_h)$ *with data [13]*

The production of the leading colorless wave packet with the desired (detected) momentum completes the process of hadronization. Any subsequent inelastic interaction is forbidden, otherwise a new hadronization process begins and the leading hadron energy falls to lower values. Such a restriction means a nuclear suppression of the production rate.

On the other hand, soft interactions of the leading quark during the hadronization in nuclear matter cannot stop or absorb the leading quark [10]. Although rescatterings of the quark in the nucleus result in an additional induced soft radiation, [10, 11], at the same time the quark



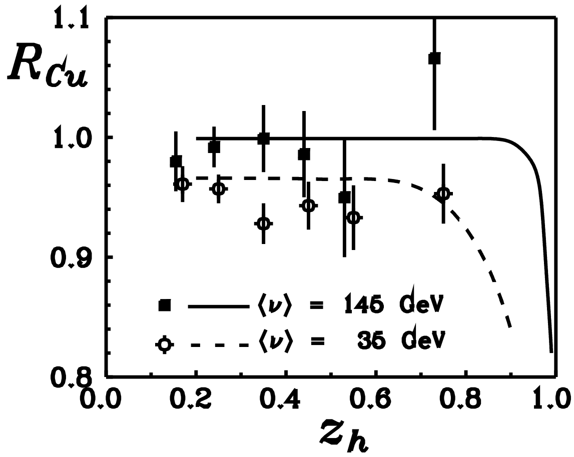
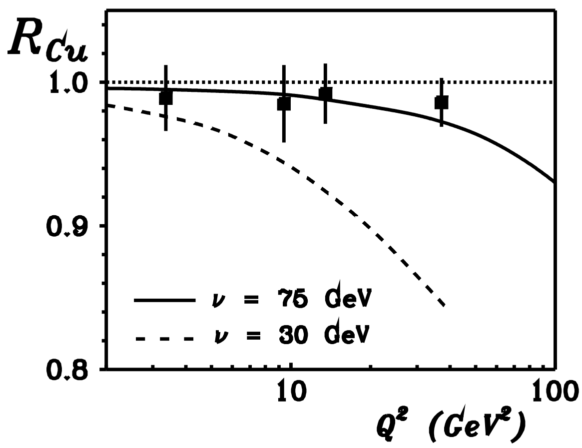


Figure 3: *Comparison of our calculations for the ν-dependence of nuclear suppression (1) integrated over z_h at $Q^2 = 6$ GeV^2 with data [13,14]*

Figure 4: *z_h-dependence of nuclear suppression.*



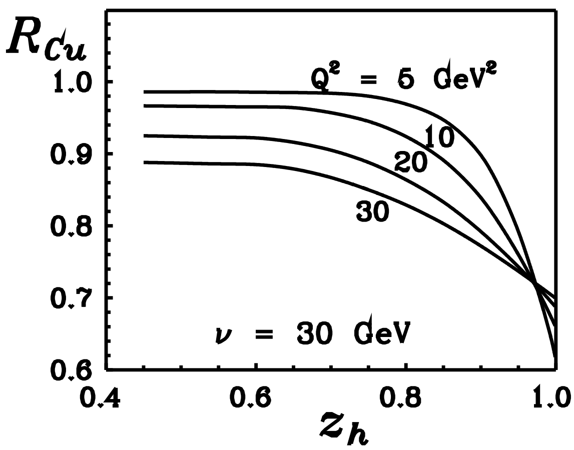


Figure 5: *Predictions for the Q^2 dependence of nuclear transparency at $\nu = 30$ and 75 GeV. Data [14] at $\nu = 75$ GeV are integrated over z_h*

Figure 6: *Predictions for the z_h dependence of nuclear transparency at $\nu = 30$ GeV vs Q^2*

looses much more energy due to the hard gluon radiation following the deep-inelastic scattering just as in vacuum. Thus, the induced soft radiation can be treated as a small correction to the energy loss and can be neglected, provided Q^2 is high enough.

The transverse size of the colorless wave packet produced in a hard reaction can be small, therefore the nuclear suppression is weaker due to color transparency. We take into account the evolution of the wave packet during its propagation through the nucleus using the path integral technique developed in [12]. Figs. 3 - 5 show quite a good agreement of our parameter-free calculations with available data [13, 9] on the ν-, z_h- and Q^2-dependence of nuclear suppression. Unfortunately, there is still no data in the region, $z_h > 0.8$.

Note that at high energies many of interesting effects go away or are difficult to observe.

Nuclear suppression integrated over z_h vanishes (see Fig. 3). The region of high z_h, where nuclear suppression is expected to be enhanced (see Fig. 4) squeezes at high ν towards $z_h = 1$, where the cross section vanishes. There is almost no Q^2-dependence at high ν and moderate values of Q^2 (see Fig. 5).

We present in Figs. 5 and 6 our predictions for the Q^2 and z_h dependence of nuclear suppression for the energy range of the HERMES experiment. We expect the onset of nuclear effects at moderate values of Q^2 (Fig. 5) as well as in the region of $z_h > 0.8$ (Fig. 6), which can be tested by the HERMES experiment.

To conclude, we have developed a phenomenology of electroproduction of leading hadrons on nuclei which is based on the perturbative QCD. Our parameter-free model is in a good agreement with available data. We stress that the energy range of the HERMES experiment is especially sensitive to the underlying dynamics of hadronization.

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