Dynamical properties of flux tubes in the Friedberg-Lee model

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Abstract

A dynamical model of confinement based on a microscopic transport description of the Friedberg-Lee model is extended to explicit color degrees of freedom. The string tension is reproduced by an adiabatic string formation from the nucleon ground state. As a particular application, we address the question of how a charmonium state might be dissociated by the strong color electric fields when moving through a color electric flux tube and speculate on the importance of such an effect with respect to the issue of $J/\psi$-suppression observed in ultrarelativistic heavy ion collisions. Furthermore, we show the dynamical breakup of flux tubes via $q-\bar{q}$-particle production and the disintegration into mesons. There we encounter some problems within the Vlasov-type realization of describing the motion of the quarks which can be resolved by a molecular dynamical approach.

1 Introduction

One of the main goals of ultrarelativistic heavy ion physics is the possible formation of a deconfined state of hot and dense nuclear matter, the quark-gluon plasma, as it is believed that the energy density reached may be high enough for its formation. On the other side, in popular microscopic models, simulating the whole reaction of the collision, strings are produced in the first few moments of the interaction, which subsequently decay (fragment) into secondaries (mesons, baryons, string-like hadronic resonances). These are best visualized as color electric flux tubes. Their creation is thought to happen by the exchange of a color octet gluon or by colorless momentum transfer among target and projectile nucleons. (The color sources in the endcaps of the flux-tubes should be thought as a quark on the one side and an anti-quark on the other side, or as a quark and a di-quark, respectively.) Before a further hadronic-like or partonic-like state of matter forms, such a temporary build-up and decay of strings

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is assumed to describe the very early collision phase. A serious problem of these models is the description of the intermediate state of the reaction, the hadronization of the ensemble of strings, which can not be treated perturbatively, since the hadrons are formed by the purely nonperturbative effects of color confinement. Therefore the dynamics of confinement should be considered in the transport descriptions.

We describe the dynamics of a flux tube within the nontopological soliton model of Friedberg and Lee [1, 2] supplemented by a semiclassical transport equation for the quarks. In this effective model the nonperturbative effects are modeled by the presence of a scalar field. Since in this model the hadron surface is generated dynamically, it is well suited for dynamical simulations. Furthermore absolute confinement is implemented by the idea of describing the vacuum as a color dielectric medium. In its original version the Lagrangian is given by

\[ \mathcal{L} = \bar{\Psi}(i\gamma_{\mu}\partial^{\mu} - m_0 - g_0 \sigma)\Psi + \frac{1}{2}(\partial_{\mu}\sigma)^2 - U(\sigma) - \frac{1}{4}\kappa(\sigma)F_{\mu\nu}^a F_{\mu\nu}^a - ig_0\bar{\Psi}g_{\mu\nu} \frac{\lambda^a}{2} \Psi A_\mu^a, \tag{1} \]

where \( \Psi \) denotes the quark fields, \( \sigma \) is the color singlet scalar field representing the long range and nonabelian effects (of multi gluon exchange). The last term contains the interaction of the residual classical and abelian color fields \( A_\mu^a \). Color confinement is obtained by the general properties of \( \kappa \). Nonabelian effects are assumed to be absorbed in the color dielectric function \( \kappa(\sigma) \) which is chosen such that \( \kappa \) vanishes as \( \sigma \) approaches its vacuum value \( \sigma = \sigma_v \) outside the bag and \( \kappa = 1 \) inside,

\[ \kappa(\sigma) = |1 - \frac{\sigma}{\sigma_{\text{vac}}}|^2 \Theta(\sigma_{\text{vac}} - \sigma). \tag{2} \]

\[ U(\sigma) = \frac{a}{2!}\sigma^2 + \frac{b}{3!}\sigma^3 + \frac{c}{4!}\sigma^4 + B \tag{3} \]

is the self-interaction potential for the scalar \( \sigma \)-field containing cubic and quartic terms.

In section II we present the resulting semiclassical transport equations [3, 4]. In section III we summarize how the parameters of the model can be fixed by static properties of hadrons and flux tubes. In section IV we adress the question what happens if a \( J/\psi \)-state moves into an environment of color electric strings and the implications of this investigation on the issue of \( J/\psi \)-suppression. Finally we present in section V the full transport dynamical approach to the disintegration of a flux tube into mesons via quark antiquark particle production. After the discussion of some intrinsic problems connected with the semiclassical limit of the model, we give the space-time description of a string-fragmentation process of the Lund-type within a molecular dynamical approach.

## 2 Transport equations

The quark degrees of freedom are effectively handled via a transport equation derived by Elze and Heinz [3] in the semiclassical limit. Since the gluon field is effectively
abelian (i.e. a Maxwell field), one can work in a U(1) sub-space of SU(3)-color. One obtains two transport equations for the phase space distributions of color charges, i.e. quarks ($f$) and antiquarks (or diquarks) ($\bar{f}$)

\[
(p_\mu \partial^\mu - m^*(\partial_\mu m^*) \partial^\mu p_\mu) f(x, p) = g_\nu p_\mu F^{\mu\nu} \partial_\nu f(x, p) \\
(p_\mu \partial^\mu - m^*(\partial_\mu m^*) \partial^\mu p_\mu) \bar{f}(x, p) = -g_\nu p_\mu F^{\mu\nu} \partial_\nu \bar{f}(x, p),
\]

(4)

which is a set of usual Vlasov equations describing the motion of charged particles in a selfconsistently generated mean scalar and vector field determined below ($m^* = m_0 + g_0 \langle \sigma \rangle$). Within this description, the coupling of quark and antiquark degrees of freedom is provided by their interaction with the mean fields $\sigma$ and the colorelectric field only. The equations are solved numerically by applying the test-particle method.

The $\sigma$ and the $A_\mu$-field are treated as classical fields (mean field approximation). The equation of motion for the scalar soliton field then reads

\[
\partial_\mu \partial^\mu \sigma + U'(\sigma) + \frac{1}{4} \kappa'(\sigma) F^{\mu\nu} F^{\mu\nu} + g_0 \bar{\Psi} \Psi = 0,
\]

(5)

which is solved using a staggered leapfrog algorithm. For the gluon field we choose the Coulomb gauge $\vec{\nabla} \cdot (\kappa \vec{A}) = 0$ resulting in

\[
\vec{\nabla} (\kappa \vec{\nabla} A_0) = -j_0,
\]

(6)

\[
-\kappa \partial_t^2 \vec{A} + \vec{\nabla}^2 \kappa \vec{A} - \vec{\nabla} \times \left( \kappa \vec{A} \times \frac{\vec{\nabla} \kappa}{\kappa} \right) = -\vec{j} + \kappa \vec{\nabla} \partial_t A_0.
\]

(7)

In the following equation (7) is neglected as it can be shown, that currents within a flux tube do not produce a magnetic field because the displacement current is exactly cancelled by the convection current if the string radius stays nearly constant (and which is exact in pure 1+1-dimensional electrodynamics). As a consequence, we determine the colorelectric field instantaneously by $\vec{E} = -\vec{\nabla} A_0$, which is determined by a two-dimensional finite element method (appropriate for cylinder symmetrical configurations).

The color charge density entering into (7) is expressed by

\[
j_0 = \frac{\eta}{(2\pi)^2} \int d^3 p (f(x, p) - \bar{f}(x, p)),
\]

(8)

whereas the scalar density $\rho_s \equiv \bar{\Psi} \Psi$ entering into eq. (5) reads

\[
\rho_s = \frac{\eta}{(2\pi)^2} \int d^3 p \frac{m^*}{\omega} (f(x, p) + \bar{f}(x, p)).
\]

(9)

(\eta \equiv 4 accounts for spin and flavour degeneracy.)

### 3 Static properties

We first examine the static limit of the transport equations. The color charge density in a hadron has to vanish locally, i.e. $j_0(x) \equiv 0$. From that one concludes that
there is no colorelectric field in the groundstate, thus $f = \bar{f}$. Because of the mean field approximation for the scalar field the distribution functions have to be of a local Thomas Fermi type

$$f(x,p) = \bar{f}(x,p) = \Theta(\mu - \omega),$$

where we have introduced the Fermi energy $\mu$ for the quarks. Considering these ground-state distributions, the quark degrees can be integrated out resulting in a selfconsistency relation for the $\sigma$-field in eq. (5). The soliton solution is constructed using a shooting method of Van Wijngaarden-Dekker-Brent [8].

The model parameters are now adjusted in a way to reproduce the quark-number, mean mass of delta and nucleon and the rms radius of the nucleon for a typical baryon. With these parameters (see table 1) we use then a new Fermi energy $\mu$ in order to go from the baryonic solution ($N_Q = 3$), resembling a quark-diquark configuration to the mesonic solution ($N_Q = 2$), resembling a quark-antiquark configuration. The physical properties of these soliton solutions are shown in table 2 and 3, where the meson mass, which typically is too high in these models, can be significantly reduced by the inclusion of momentum projection methods and the colormagnetic energies [2]. The remaining properties are very well reproduced within our model approach. These groundstate solutions will be used in the following sections for dynamical applications.

| parameter set          |            |
|------------------------|------------|
| $a$ [fm$^{-2}$]        | 0.0        |
| $b$ [fm$^{-1}$]        | -419.3     |
| $c$ [1]                | 4973.0     |
| $B$ [fm$^{-4}$]        | 0.283      |
| $g_0$ [1]              | 8.0        |
| $m_0$ [fm$^{-1}$]      | 0.025      |
| $\mu$ [fm$^{-1}$]      | 1.768 / 1.9582  |
| $\eta$                | 4          |
| $\sigma_{vac}$ [fm$^{-1}$] | 0.253    |
| glueball mass [GeV]    | 1.045      |

Table 1: The table shows the parameters used to construct a mesonic and a baryonic soliton solution of the sigma field equation. The first value of $\mu$ is used for the baryonic solution, the second one for the mesonic.

Since the colorelectric field vanishes in the static solution, the strong coupling constant $\alpha_s = g_0^2/4\pi$ cannot be fixed from ground state properties. In order to determine its value we investigate the colorelectric energy in the cavity by forming a string like configuration (see fig. 1). We find the typical string like behaviour of the $\vec{E}$-field [3], namely being constant and almost independent of the axial coordinate in the region between the quarks. The genuine feature of this field is, that it is parallel everywhere.
### Parameter set

| Parameter set | meson | Experimental Data |
|---------------|-------|-------------------|
| $E \ [MeV]$  | 798   | 465               |
| $RMS \ [fm]$ | 0.653 | 0.66              |

Table 2: The table shows the results of the fits for the meson compared to the experimental data. The experimental data are the mean values of the pion and the rho-meson (taken from [2]).

| Parameter set | baryon | Experimental Data |
|---------------|--------|-------------------|
| $E \ [MeV]$  | 1099   | 1087              |
| $RMS \ [fm]$ | 0.693  | 0.83              |

Table 3: The table shows the results of the fit for the baryon compared to the experimental data. The experimental data are taken from [2].

at the boundary of the cavity, which is due to the von Neumann boundary conditions, forcing the normal component of the displacement field $\vec{D} = \kappa \vec{E}$ to vanish. Determining the colorelectric energy $E_{\text{glue}} = \frac{1}{2} \int d^3r \vec{D} \vec{E}$ and the energy carried by the scalar field as a function of the length of the string provides us with a measure for the string constant $\tau$. In order to obtain the generally accepted value of $\tau \approx 1 \text{GeV/fm}$, we need to have values of $\alpha_s \approx 2$, in accordance with other estimates [2]. The volume averaged electric field $\langle E_z \rangle$ takes a value of $1.6 \text{ fm}^{-2}$ with $E_{z,\text{max}} = 2.2 \text{ fm}^{-2}$ along the $z$-axis. The radius of the string results as approximately $0.8 \text{ fm}$.

One can also extract a string-string interaction potential by an adiabatic fusion of two parallel or antiparallel strings [4]. The potential is obtained by comparing the colorelectric and the $\sigma$–field energies of the fused configuration to the (infinitely) separated configuration. Due to confinement in this model there is no long range interaction (over distances typically of the order of $1 \text{ fm}$ and more). In the case of two parallel strings, in the intermediate range, when the confinement wall starts to break down, we have a balance of gaining surface energy from the $\sigma$–field and increasing the total energy from the coherent addition of the colorelectric field. In the short range region the potential is strongly repulsive due to the unscreened quark color charges at the endcaps.

4 \textit{J}/\psi\textit{-dissociation by a color flux tube}

Since the original work of Matsui and Satz [9] who proposed $J/\psi$–suppression as a clear signal for quark gluon plasma formation (the electric forces responsible for the binding are sufficiently screened by the plasma so that no boundstate should form) an intense work by several experimental groups (NA38-collaboration) has addressed this issue [10].
Indeed a suppression in the corresponding dilepton signal compared to the Drell-Yan background has been seen over the last years for lighter projectiles like O+Cu, O+U and S+U. However, these observations may also be explained alternatively: (1) A part or even the total suppression can probably be explained by $J/\psi$-absorption on the surrounding nucleons [11], i.e. $J/\psi + N \rightarrow \Lambda_C + \bar{D}$; (2) additional absorption by exothermic reactions like i.e. $J/\psi + \rho \rightarrow D + \bar{D}$ might be attributed to ‘comovers’ (‘mesons’) being produced as secondaries. Indeed, if $\sigma_{\psi N} \approx 7$ mb is assumed, the reported suppression could be nicely reproduced by the absorption within the nuclear environment. In any case, the general consensus has been, that a highly (energy) dense intermediate reaction zone is needed to explain the observed suppression. Recently the data taken in Pb+Pb reactions at 160 AGeV by the NA50-collaboration [12] have given new excitement to the whole issue as a stronger absorption has been found than suggested by the models of absorption models on nucleons.

The $J/\psi$ (or better $c\bar{c}$-state which later will form the $J/\psi$) as a rather heavy hadronic particle is being produced at the earliest state of the reaction in a hard collision among the nucleons. Thus it is natural to ask what happens if a (pre-) $J/\psi$-state (as a strongly interacting probe) enters into an environment of color electric strings. As the string carries a lot of internal energy (to produce the later secondaries) the quarkonia state might get absorbed and completely dissociated by the intense color electric field inside a single flux tube.

To address this question, we first describe semiclassically a bound $c\bar{c}$-state (synonymously a $J/\psi$) and also a bound $c\bar{q}$-state (synonymously a $D$) within our model. We start from a light meson state and then adiabatically increase the mass of the initial light quark(s) ($m_{i,q} \approx 10$ MeV) to a final heavy quark mass (for details see [13]). Taking the a priori unknown charm quark mass as $\approx 1.25$ GeV, the overall mass of the states are found as

$$m_{c\bar{c}} \approx 3 \text{ GeV}$$
$$m_{c\bar{q}} \approx 1.7 \text{ GeV}.$$  \hspace{1cm} (11)

The mass of the $c\bar{c}$-state lies only slightly above the experimental value by about $\approx 60$ MeV, whereas the mass of the $c\bar{q}$-state is lowered by about the same value, thus providing a very good description. The radius of the heavy meson states (the ‘$J/\psi$’- and ‘$D$’-meson) are in the right range of what one would expect from other quarkonia model.

We now insert the ‘$J/\psi$’ ‘by hand’ into the central interior of a stationary flux tube (see figs. 2). Initializing a cylindrically symmetric configuration in this way, we proceed by solving the full set of dynamical equations of motion, i.e. the two Vlasov equations (4) for the heavy quark and antiquark distribution, the equation (5) for the soliton field $\sigma$ and the one for the electric potential (6). What happens is that the heavy quark is steadily pulled to one endcap of the string carrying opposite charge, whereas the antiquark correspondingly moves into the other direction. Being pulled apart, the two
displaced charges screen the overall electric field in between. In other words, the electric field energy originally stored is transformed into the kinetic energy of the oppositely moving two heavy quarks. In figs. 2 also the evolution of the $\sigma$-field is depicted for various time stages. At the initial times one clearly sees how the soliton solution of the flux tube is distorted by introducing the $c\bar{c}$-state into the center of the string. At the timesteps up to about 5 fm/c the $\sigma$-field follows the quark distribution quite well, its shape reflects the shape of the latter distribution. In the further evolution (not shown) the soliton field relaxes much more slowly to its vacuum value compared to the more or less instantaneous screening of the electric field. It typically takes about 5 fm/c until $\sigma$ turns for the first time to its vacuum value. This long time interval is basically a consequence of the nonlinear potential $U(\sigma)$ being flat around $\sigma \approx 0$ and corresponds to the time needed for the soliton field to ‘move down’ the potential hill to its vacuum value at $\sigma = \sigma_V$. The original field energy carried by the soliton field (the ‘Bag’ energy) being still present after the heavy quarks have already moved apart is initially more or less conserved. The late oscillations observed in the simulations are only damped by a transverse expansion of the field itself, which, however, is a rather slow process because of the large ‘glueball’ mass of the field. If one would consider a chiral $O(4)$-extension for the soliton field one expects that these oscillations might accordingly transform into low-momentum (and nearly massless) pion modes and should accordingly be damped away more quickly.

From this simulation we conclude that a localized $c\bar{c}$-state immediately gets separated by the strong colorelectric field, and, in return, will finally be disintegrated into $D$-meson (or $\Lambda_C$-baryon) like configurations. The dissociation corresponds to classical field ionization. Such a dissociation should also persist if initially the $c\bar{c}$-state enters the string with some moderate or even quite large transverse momentum [13].

But what is now the actual situation in a relativistic heavy ion collision? If one believes in the successful descriptions of microscopic approaches, a large region in space in the first few moments after the reaction is spanned by highly excited longitudinal strings. Especially for the more heavy systems the effective region (or volume) of all the strings being produced within a short time interval of less then one fm/c gets so large that the strings become closely packed or already overlap. Although most often the strings are thought to fragment independently one might also consider the possibility of color rope formation as higher charged tubes. For a quantitative estimate we depict in fig. 3 the position of nucleonic strings (highly energetic strings excited by a target and projectile nucleon) all being produced in a cms-time interval of $\approx 0.5$ fm/c within a slab of the complete transverse area and 1 fm thickness along the longitudinal direction in the cms frame in a central S+U reaction at 200AGeV generated within the HSD approach [14]. This simple reasoning illustrates that the whole reaction area in an ultrarelativistic heavy ion collisions might be completely filled by strings in the first few moments. If such a large spatial region does exist, most of the produced $J/\psi$’s have to pass it and thus are affected by this highly excited environment. The
suppression of the $J/\psi$’s will depend crucially on the (average) length of the produced strings before they hadronize, as this decides about the available ‘empty’ space for the $J/\psi$’s (or pre-$J/\psi$-states) to escape the initial reaction zone without entering any of the individual strings being built up. Such a conclusion can only be tested further and strengthened if one incorporates these ideas within one of the present microscopic transport algorithms.

5 Flux tube breaking

In this section we show how a flux tube of the Friedberg-Lee model breaks up due to quark-antiquark pair production. One usually describes such processes with the Schwinger formula, which gives a constant pair production rate of electron-positron pairs in QED, depending only on the absolute value of the electric field. However, this formula cannot be naively transferred to the QCD case of $q\bar{q}$–production in a flux tube, because the back reaction of the produced pairs on the external field has to be considered, since this screening of the field is finally responsible for the breakup of the strings. In addition one must consider the modifications of the pair production rate due to transverse and longitudinal confinement. A detailed analysis of the time dependent pair production process within the Friedberg-Lee model has not yet been performed. Therefore we have to rely on different phenomenological arguments, providing us with a simple guideline to the space-time evolution of a fragmenting string. The earliest and most successful model of string fragmentation is the Lund-model [15]. Within this model one assumes the quarks to be massless and therefore moving with the speed of light. Even the original $Q\bar{Q}$–pair, generating the string, is supposed to travel on the lightcone. Within these assumptions it is sufficient to describe the fragmentation process by a 1+1-dimensional space-time geometry, for it is causally impossible that the produced $q\bar{q}$–pairs propagate towards the endcaps (the $Q\bar{Q}$–pair).

Hence, as a prerequisite we assume that the original $Q\bar{Q}$–pair travels with the speed of light along the $\pm z$–axis, irrespective of the interior dynamics, which is generated by solving selfconsistently the transport equations (1) together with the mean-field equations for the $\sigma$–field (3) and the colorelectric field (4). The produced $q\bar{q}$–pairs are inserted into the dynamically evolving flux-tube by assuming their groundstate spatial shape that has been determined in section II and with a vanishing total momentum [6]. If one inserts only one (light) $q\bar{q}$–pair, the fragmentaion of the strings follow directly to what one would expect from the Lund phenomenology [4]. The inserted two light quarks get immediately dissociated and do follow the two endcaps propagating along the lightcone. The colorelectric field in between gets completely screened, whereas the scalar fields is a little retarded in its motion. From the left part of fig.4 we recognize a small dispersion of the $q$ and $\bar{q}$ distributions at late times, which is caused by the selfinteraction inside of the quark distributions, leading to a screening of the charges: a charge fragment at the end of one of the two substrings feels a smaller colorelectric
force due to its interaction with the charge fragments in front of it. This is clearly seen in the right part of fig. 4 where the longitudinal momentum distribution of the testparticles describing the propagation of the two inserted quarks is shown. The momentum distribution in $p_z$ of one light quark is spread from nearly 0 up to 5 GeV/c!

As a further example we show in fig. 5 (left part) the space-time evolution of the string fragmentation via equal time production of three $q\bar{q}$–pairs within the full model. In the upper row at $t = 2.8\, fm/c$, the string is already extended to a length of 5.6 fm; the constant colorelectric field along the string is reflected in the constant field energy between the charges. At $t = 3.0\, fm/c$ the three $q\bar{q}$–pairs are inserted at $z = 0\, fm$ and $z = \pm 2.0\, fm$ with a vanishing relative velocity. In the following time evolution we observe, as in the previous example, the formation of the outer two meson pairs propagating on the lightcone. In the inner region of the flux-tube, the respective quark pairs penetrate each other, having already acquired a small dispersion. At the time $t = 4.2\, fm/c$, the colorelectric field is almost completely screened. This abrupt change of the colorelectric field leaves the $\sigma$–field highly excited. Therefore the $\sigma$–field snaps off and starts to oscillate around the nonperturbative vacuum value $\sigma_{vac}$. Thus the $\sigma$–field, which is supposed to localize the quarks, cannot prevent a further dispersion of the color charges and the latter neutralize and dissolve along the $z$-axis to a length of about $8\, fm$ at $t = 8.2\, fm/c$. We are thus faced with the fact of having produced real quark-antiquark fragments evaporating in the nonperturbative QCD vacuum.

As already mentioned, in the present numerical realization of our model each test-particle of the respective ensembles interacts not only with the testparticles of the other quarks, but also with the testparticles of the same quark. This selfinteraction, which leads to the dispersion of the quark distributions (see e.g. fig.4), becomes dominant in the limit of only a few interacting physical particles; in the extreme case, a single color charge distribution with an almost vanishing mass would immediately dissolve due to the selfinteraction. Unlike in other models, like e.g. the Walecka-model, which also contains an attractive scalar field and a repulsive vector field and which has been treated successfully, in the present case the attractive $\sigma$–field acts only in transverse direction on the produced quarks. Its longitudinal effect is completely suppressed by the colorelectric field. Thus we are forced to formulate a new dynamical approach, that conserves the identity of the different quarks and antiquarks and is free of the selfinteraction problem: A possible way to overcome the problems mentioned above is to treat the quarks in a molecular dynamical framework [4]. In this approach each physical particle (quark or antiquark) is simulated by just one particle and not by an ensemble of testparticles. Due to the fact that point-like particles cause short range divergences in the field equations, a specific spherically symmetric distribution is assigned to each individual testparticle. A testparticle then describes the motion of the center of its distribution according to the Hamiltonian equation of motion.

In the right part of fig. 5 we show the corresponding result within this molecular dynamical approach. We see how the screening proceeds not only at the coordinate
center, but also between the other fragments. In the third and fourth row we clearly see how the two lightcone fragments have been formed as well as the two excited 'Yo-Yo'-modes. The confining $\sigma$-field follows the motion of the quarks with a certain retardation, but finally settles into individual stable soliton configurations. In summary, we find a more or less complete correspondence of the dynamical behaviour to the one expected from 1+1-dimensional Lund phenomenology.

6 Summary and Outlook

We have given our present status of a dynamical realization of the Friedberg-Lee model. This model describes absolute confinement phenomenologically. The transport equations for the quarks are of Vlasov-type, whereas the soliton field and the chromoelectric field are treated as (classical) mean fields. The static properties of meson and baryons as well as the description of color electric flux tubes can be successfully be described within this approach.

As a first dynamical application we have investigated how a $c\bar{c}$-state (a $J/\psi$) behaves inside a chromoelectric flux tube. Due to the strong color electric field inside the flux tube the heavy meson becomes dissociated (due to the complete screening in effective 1+1-dimensions) by field ionization rather immediately (on a timescale of $\approx 1 \text{ fm/c}$) ending up finally in $D$- or $\Lambda_c$-like states. It is expected that a large region in space in the first few moments in an ultrarelativistic heavy ion collision is spanned by highly excited longitudinal strings, especially for the more heavy systems. We have speculated that the dissociation process might provide an alternative or additional source for explaining the observed $J/\psi-$suppression. This possible absorption mechanism will be implemented within one of the present microscopic transport algorithm in order to verify these ideas.

Secondly we have shown how the flux-tube breaking proceeds in the Friedberg-Lee model. In the case of three produced $q\bar{q}$-pairs, the outer meson pairs are formed, but in the center of the fragmenting string, where the two parallel 'Yo-Yo'-modes are supposed to build up, we observe a different behaviour: When the respective quark distributions start to penetrate and neutralize, a small dispersion of these is already observed. At this time the colorelectric field is completely screened in the center of the string. This sudden change of the source of the flux-tube leaves the $\sigma$-field highly excited and thus it starts to oscillate around the nonperturbative vacuum value. The $\sigma$-field is not able to prevent the further dispersion of the quark distributions, which neutralize and dissolve along the string-axis. With the scalar density being significantly reduced in this case, the $\sigma$-field undergoes a complete phase transition to the nonperturbative vacuum.

On the other hand, the dispersion of the quark distributions is caused by their selfinteraction. This dispersion is caused by the fact, that we deal with only a few quarks that are treated as classical charge distributions and interacting strongly via a
classical confining potential. In this case the selfinteraction of the color charges becomes dominant and generates a screening of the charges, which finally is responsible for the dispersion of these. We have shown a possible way to overcome these problems by treating the quarks and antiquarks as point particles with an appropriately chosen charge distribution as commonly used in molecular dynamical simulations.

Summarizing we can say that our model provides us with a useful tool to describe the full dynamical evolution of string formation and decay via multiple quark-antiquark production. The aim for future investigations should be to apply our model to more realistic scenarios, like multiple string production and decay in relativistic heavy-ion collisions, for describing the formation of hadrons out of an expanding and cooling quark-gluon plasma.

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Figure 1: The electric field $E$ inside a cavity (flux tube) of approximately 2 fm and 12 fm in length, respectively.

Figure 2: The quark distribution (left side) and the scalar field distribution (right side) at various times in the evolution ($t = -0.2, 0.8, 3.4, 4.6 \text{ fm/c}$) are shown in a contourplot.

Figure 3: The position of nucleonic strings is plotted in the cms frame in a central S+U reaction at 200AGeV. Each string is accompanied by a circle of radius 0.5 fm.

Figure 4: The left figure shows the color charge density $\rho$ at a late stage ($\approx 5 \text{ fm/c}$) for the breakup of the a string via the production of one $q\bar{q}$ pair at time $t = 0 \text{ fm/c}$. The right figure depicts the longitudinal momentum distribution of the testparticles of the inserted two light quarks at that time.

Figure 5: The figure shows the time evolution of the color charge density for the breakup of the string via the production of three $q\bar{q}$ pairs at different time steps $t = 2.8, 3.6, 4.2, 8.2 \text{ fm/c}$. In the left part the motion of the quarks has been calculated within the full transport dynamics, whereas in the right part the molecular dynamical approach has been used. The equidistant contour lines run from $-0.5 \text{ fm}^{-3}$ to $0.5 \text{ fm}^{-3}$.
