$J/\psi$–dissociation by a color electric flux tube

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Abstract

We address the question of how a $c-\bar{c}$-state (a $J/\psi$) can be dissociated by the strong color electric fields when moving through a color electric flux tube. The color electric flux tube and the dissociation of the heavy quarkonia state are both described within the Friedberg-Lee color dielectric model. We speculate on the importance of such an effect with respect to the observed $J/\psi$–suppression in ultrarelativistic heavy ion collisions.

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I. INTRODUCTION AND MOTIVATION

Since the original work of Matsui and Satz [1] 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 [2]. 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 [3], i.e. \( J/\psi + N \rightarrow \Lambda_c + D \); (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 [4]. Indeed, if \( \sigma_{\psi N}^{\text{abs}} \approx 7 \text{ mb} \) is assumed, the reported suppression could be nicely reproduced by the absorption within the nuclear environment. In any case, the general consensus has then been that a highly (energy) dense intermediate reaction zone was needed to explain the observed suppression.

Recently the data taken in Pb+Pb reactions at 160 AGeV by the NA50-collaboration [5] 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. It was immediately speculated [6] that a large inner region of the reaction zone must absorb all \( J/\psi \) (being ‘black’) in order to be compatible with observation. Though it was reported by Gavin and Vogt [7] and by Cassing and Ko [8] that the comover model also could explain the stronger suppression as relatively more mesons are produced for large \( E_t \) events (or fluctuations) than for lighter systems, the latter calculations also show that this suppression takes place at energy densities that are much higher than those in which normal hadrons are expected to survive.

In popular microscopic models, simulating the whole reaction of the collision, in the first few moments interaction strings are produced 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 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.) These strings might also overlap (if their density is large) to form higher charged tubes, so-called color ropes \[9\]. In any case, before a further hadronic-like or partonic-like state of matter does form, such a temporary build up and decay of strings is assumed to describe the very early collision phase. In fact one might call such an environment a precursor of quark matter.

The \(J/\psi\) (or better \(c\bar{c}\)-state which later will form the \(J/\psi\)) as a rather heavy hadronic particle is also 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 moves (or 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. This intuitive idea is the basic reason for our work.

The dynamics of the flux tube we describe within the Friedberg-Lee model \[10,11\] by a semi-classical transport algorithm (section II). As a first step, within this model, we generate a bound \(c\bar{c}\)-state with reasonable properties in order to mimic a \(J/\psi\) (section III). We then adress how such a state gets dissociated when inserted into a chromoelectric flux tube and how the tube subsequently does split (section IV). We conclude with a summary and the implications of our investigation on the issue of \(J/\psi\)-suppression.

II. DYNAMICAL REALIZATION OF THE FRIEDBERG-LEE MODEL

Recently we have written down semiclassical transport equations for quark distributions \[12,13\] starting from the phenomenological Lagrangian of Friedberg and Lee \[10,11\]

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

where \(\Psi\) denotes the quark fields, \(\sigma\) is the color singlet scalar field representing the long range and nonabelian effects (of multi gluon exchange), and the last term contains the interaction of
the residual classical and abelian color fields $A_\mu^a$. All the 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,

$$\kappa(\sigma) = |1 - \frac{\sigma}{\sigma_{vac}}|^2 \Theta(\sigma_{vac} - \sigma). \quad (2)$$

$$U(\sigma) = \frac{a}{2!} \sigma^2 + \frac{b}{3!} \sigma^3 + \frac{c}{4!} \sigma^4 + B \quad (3)$$

is the self-interaction potential for the scalar $\sigma$-field containing cubic and quartic terms. Details of the numerical realization and the further approximations made are found in [13,14]. One obtains two transport equations for the phase space distributions of quarks ($f$) and antiquarks ($\bar{f}$)

$$\left(p_\mu \partial^\mu - m^*(\partial_\mu m^*) \partial^\mu_p\right)f(x,p) = g_v p_\mu F^{\mu\nu}_c \partial^\nu_p f(x,p)$$

$$\left(p_\mu \partial^\mu - m^*(\partial_\mu m^*) \partial^\mu_p\right)\bar{f}(x,p) = -g_v p_\mu F^{\mu\nu}_c \partial^\nu_p \bar{f}(x,p), \quad (4)$$

which are numerically solved by applying the test-particle method, a classical equation of motion for the scalar soliton field

$$\partial_\mu \partial^\mu \sigma + U'(\sigma) + \frac{1}{4} \kappa'(\sigma) F^c_{\mu\nu} F^{\mu\nu}_c + g_0 \bar{\Psi} \Psi = 0, \quad (5)$$

which is solved using a staggered leapfrog algorithm [12], and an equation of motion for the confined colourelectric field $\vec{D} = \kappa \vec{E}$

$$\vec{\nabla} (\kappa \vec{\nabla} A_0) = -j_0, \quad (6)$$

which is determined by a two-dimensional finite element method (appropriate for cylinder symmetrical configurations) developed by Mitchell [15].

First we note that here we explicitly suppress the color magnetic fields, i.e. $F^c_{\mu\nu} F^{\mu\nu}_c \rightarrow 2\vec{E}^2$. (In a string like configuration with nearly constant radius Wilets and Puff have shown that due to the displacement current in the Maxwell equation magnetic fields are not produced [17].) $f$ and $\bar{f}$ describe the distribution of charged quarks and anticharged antiquarks (or diquarks) so that the colour charge density is expressed by
\[ j_0 = \frac{\eta}{(2\pi)^3} \int d^3p (f(x,p) - \bar{f}(x,p)) , \] (7)

whereas the scalar density \( \rho_s \equiv \bar{\Psi} \Psi \) entering into eq. (4) reads
\[ \rho_s = \frac{\eta}{(2\pi)^3} \int d^3p \frac{m^*}{\omega} (f(x,p) + \bar{f}(x,p)) . \] (8)

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

Within this approach one succeeds in simulating dynamically confinement as (a) a single charged quark is forced back into the configuration by the dielectric displacement \( \vec{D} \) if it tries to escape [11]; and (b) the dielectric displacement \( \vec{D} = \kappa \vec{E} \) stays also within the configuration as it has to vanish in the outside (and nonperturbative) vacuum where \( \kappa \to 0 \) [16].

Static properties of soliton-like nucleons and mesons are successfully reproduced by an appropriate fixing of the scalar potential \( U(\sigma) \) and the coupling constant \( g_s \). The vector coupling \( \alpha_s \equiv \frac{g^2}{4\pi} = 1.92 \) is fixed by assuming an effective string constant \( \tau = 1 \) GeV/fm when simulating a colourelectric flux tube [13]. Here the volume averaged electric field \( \langle E_z \rangle \) takes a value of 1.6 fm\(^{-2}\) with \( E_{z,max} = 2.2 \) fm\(^{-2}\) along the z-axis. The radius of the string results as approximately \( 0.8 \) fm.

**III. J/ψ- AND D,\bar{D}-STATES**

As a first step, within our model, we want to describe semiclassically a bound \( c\bar{c} \)-state (synonymously a \( J/\psi \)) and also a bound \( c\bar{q} \)-state (synonymously a \( D \)). We start from a light meson state described within our model (using the parameters of the model [12–14]) and then adiabatically increase the mass of the initial light quark(s) \( m_{i,q} \approx 10 \) MeV) to a final heavy quark mass \( m'_{f,c} = 1.56 \) GeV (taken from ref. [1]), i.e.

\[ m_q(t) = m_{i,q}(t = 0) + (m_{f,c} - m_{i,q}(t = 0)) t/\tau_{s.o.} \text{ for } 0 < t \leq \tau_{s.o.} \] 
\[ m_q(t) = m_{f,c} \text{ for } t > \tau_{s.o.} . \] (9)

The change of the meson state with mass of both (or one) quarks is thus described fully dynamically. In figs. 1 we see how the bag size shrinks from \( r_{cms} \approx 0.66 \) fm to \( r_{cms} \approx 0.25 \) fm for the case
of a $c\bar{c}$-groundstate and to $r_{\text{cms}} \approx 0.40$ fm for the case of a $c\bar{q}$-groundstate for a switching on time of $\tau_{\text{s.o.}} = 20$ fm/c. We checked our results for various switching on times $\tau_{\text{s.o.}} = 20, 30, 50$ fm/c and found that the state described does not depend on the parameter $\tau_{\text{s.o.}}$. The shrinking of the states is in accordance with the expectation of any baglike model. When the quark masses are increased to their final value, the pressure the more heavy quarks assert on the surface drops too, thus leading to the shrinking of the whole bag until it stabilizes at a smaller radius where pressure equilibrium is reached again. 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. In addition, fig. 2 shows the spatial quark distribution and the soliton solution $\sigma(r)$ as a function of the radius for the $c\bar{c}$-state at the beginning and the end of the evolution according to (9).

When determining the overall mass of the states we find for $m_c = 1.56$ GeV

$$m_{c\bar{c}} = 4.10 \text{ GeV}$$
$$m_{c\bar{q}} = 2.25 \text{ GeV}.$$  

These numbers have to be compared to (a) the mass of the $\eta_c$: $m_{\eta_c} = 2.94$ GeV and the $J/\psi$: $m_{J/\psi} = 3.097$ GeV and (2) to the mass of the $D$-meson: $m_D = 1.89$ GeV and the $D^*$: $m_{D^*} = 2.01$ GeV. The $c\bar{c}$-state generated within the model (and within the parameters chosen to describe the light mesons) are thus about 1 GeV too heavy, the $c\bar{q}$-state about 0.4 GeV, whereas the difference $2m_{c\bar{q}} - m_{c\bar{c}} \approx 0.4$ GeV is in quite good agreement to $2m_D - m_{J/\psi} = 0.68$ GeV. This discrepancy one can relax by choosing a smaller bare mass of the charm quark. This mass is rigorously not known and only extracted by various charmonium models and might vary between 1 to 1.6 GeV [18]. Taking the a priori unknown charm quark mass as $\approx 1.25$ GeV (compare figs. 1), the mass of the $c\bar{c}$-state would lie slightly above the experimental value by about $\approx 200$ MeV, whereas the mass of the $c\bar{q}$-state is lowered by about the same value, thus providing a much better description within our model. In the following we stay within our model states (10) (described by a bare mass of $m_c = 1.56$ GeV) as our conclusions will not depend on their masses and the results given in the next section will only weakly be affected by assuming a different mass for $m_c$ within the afore mentioned range.
We now turn to the main issue of the present work: What happens to a $c\bar{c}$-state entering the region spanned by a single chromoelectric flux tube?

We prepare a chromodielectric flux tube (see also [13]) by pulling the opposite charge distributions of a (light) groundstate meson steadily apart. The quarks are propagated with a constant velocity up to a final spatial extent of about 8 fm (see e.g. fig. 3) within a time interval of about 13 fm/c. At that given time $t_0 \approx 13$ fm/c we keep the quark distributions at the endcaps fixed, thus providing the sources for the chromodielectric field. We then insert the ‘$J/\psi$’ ‘by hand’ into the central interior of the (by now) stationary flux tube, where the initial momentum and spatial distributions of the heavy quarks are chosen as the one obtained in the previous section. Initializing the 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).

In fig. 3 we show the evolution of the total electric field energy as a function of time. As one notices, after inserting the bound state, the field energy decreases approximately linearly with time. What happens is that the heavy quark (described by its phase space distribution) is steadily pulled to one endcap of the string carrying opposite charge, whereas the antiquark carrying the opposite charge is correspondingly pulled into the other direction (see figs. 4). 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. Being accelerated, after some smaller time interval $\Delta t \approx 1$ fm/c, both heavy quarks have nearly reached the speed of light, so that the linear decrease of the overall electric field energy depicted in fig. 3 becomes obvious.

In figs. 4 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 18 fm/c the \(\sigma\)-field follows the quark distribution quite well, its shape reflects the shape of the latter distribution. At 20.8 fm/c the inner heavy charges have reached the fixed endpoints of the flux tube. At this point we stop – for numerical reasons – the further propagation of the quarks and concentrate on the further time development of the \(\sigma\)-field shown in figs. 5. In this further evolution 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 \([14]\). 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 \(\sigma = \sigma_V\). Even when \(\sigma_V\) has been reached, though, the field still starts to oscillate around this value which can be seen in figs. 5: In the second row, at \(t=23.2\) fm/c, the \(\sigma\)-field value is slightly above \(\sigma_{\text{vac}}\), whereas in the third row, at \(t=26.8\) fm/c, its value is nearly in the perturbative region. 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. Thus it takes up to \(\approx 15\) fm/c until the \(\sigma\)-field has finally relaxed and the static heavy \(D/\bar{D}\)-meson states are formed. 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 our simulation we conclude that a localized \(c\bar{c}\)-state immediately gets separated by the strong colourelectric field, and, in return, will finally be dissociated into \(D\)-meson (or \(\Lambda_C\)-baryon) like configurations. One wonders if such a dissociation would persist also if initially the \(c\bar{c}\)-state enters the string with some moderate transverse momentum. The finite element method implemented in order to calculate the chromoelectric potential \(A_0(r, z)\) can presently only handle cylindrically symmetric configurations (as e.g. a static string). Hence we are yet not able to simulate the full dynamics when a \(c\bar{c}\)-state (as generated in the previous section) enters a static and elementary flux tube with some transverse momenta and finite impact (though this work is
in preparation [20]). In the following we give an estimate that such a disintegration of a $c\bar{c}$-state should still persist for even quite large transverse momenta.

The time needed to pass the whole flux tube is roughly estimated as

$$\Delta t \approx \frac{2r_{\text{rms}}}{v_{t/\psi}} \approx 2r_{\text{rms}} \frac{\sqrt{p_t^2 + m_{J/\psi}^2}}{|p_t|} \approx 2 - 4 \text{ fm/c}$$

(11)

for moderate to large (transverse) momenta in the range from 1 to 5 GeV/c. For an average field strength $\langle |E_z| \rangle$ of 1.61/fm$^2$ the heavy quarks will suffer a gain in longitudinal momentum given approximately by ($\alpha_s = g_v^2/4\pi \approx 2$)

$$|\Delta p_z| \approx g_v |E_z| \Delta t \approx 1.6 \text{ GeV/fm} \Delta t \approx 3 - 6 \text{ GeV/c} ,$$

(12)

whereas its internal rise in longitudinal kinetic energy of both quarks (by assuming that each heavy quark moves, on the average, with, let us say, half speed of light)

$$\Delta W \approx 2g_v |E_z| \Delta s \approx g_v |E_z| \Delta t \approx 3 - 6 \text{ GeV} .$$

(13)

Such an increase in momenta or internal kinetic energy is certainly sufficient to break up the $c\bar{c}$-state into two D-meson like states, as the energy needed is given by the difference in mass, i.e. of about 0.68 GeV. As the strings are thought to be energetic enough to produce all secondaries over the whole rapidity range in the heavy ion collision, its internal color electric field should be large enough to break up even more heavily bound systems like the $J/\psi$. We are hence quite confident that our observed dissociation would also be found within a full 3-dimensional simulation of a $c\bar{c}$-state interpenetrating a chromodielectric flux tube with a nonvanishing transverse momentum.

In some popular charmonium models a rather strong effective Coulomb potential between the $c$ and $\bar{c}$ quark of the form $V_{c\bar{c}} \sim -C/r$ , to be valid at short distances, gives a dominant contribution to the overall spectra and gives rise to a binding energy estimated to be 800 MeV for the $J/\psi$ as a 1s-state. When modeling this state within our semiclassical approach, the groundstate is only bound by the strong confining scalar field. The dissociation described above corresponds to classical field ionization. The same effects exist in a quantum mechanical treatment: A quantum
mechanical system of two heavy quarks will dissociate at long times if a constant background field with a potential \( V \sim -g_v Ez \) is switched on \(^{21}\). For our specific numbers, along the z-direction, this leads to a potential of \( V \approx (-1.7)z \text{ GeV/fm} - (-2.1)z \text{ GeV/fm} \) due to the strong electric field. The modifications are thus strong and nonperturbative. For \( z = 0.3 \text{ fm} \) the effective total potential \( V_{\text{eff}}(z) \sim -C/z - g_v Ez \) is already lowered by \( \approx 600 \text{ MeV} \) compared to the ‘bare’ potential \(-C/z\). The existence of a (then highly modified) resonant 1s-state seems very unlikely.

In addition one also can consider this quantum mechanical problem as a nonstationary situation when the \( J/\psi \) passes through the tube with some transverse velocity. When entering, the abrupt change in situation can not be considered as adiabatic, but better as diabatic: The original bound state has no time to adiabatically lower its energy by at least 0.6 GeV, but will immediately dissolve into the continuum.

V. CONCLUSION AND OUTLOOK

In this work we have adressed the question of how a \( c\bar{c} \)-state (a \( J/\psi \)) behaves inside a chromoelectric flux tube. The flux tube as well as the \( c\bar{c} \)-groundstate are described within the model by Friedberg and Lee \(^{11}\). Due to the strong color electric field inside the flux tube the heavy meson becomes dissociated rather immediately (on a timescale of \( \lesssim 1 \text{ fm/c} \)) ending up finally in \( D \)- or \( \Lambda_c \)-like states. We could not simulate a fully 3-dimensional situation when the heavy meson enters the string with some fixed transverse momentum because our present numerical realization of the model works only for cylindrically symmetric configurations. However, we give an estimate that such a dissociation should also occur by field ionization for moderate and higher transverse momenta of the \( c\bar{c} \)-state when passing through the tube within our semiclassical approach.

We believe that this investigation adresses an intriguing question for understanding particular topics of QCD with respect to relativistic heavy ion collisions. If one believes in the succesful descriptions of microscopic approaches (like FRITIOF \(^{22}\), RQMD \(^{23}\), VENUS \(^{24}\), HSD \(^{25}\)), a large region in space in the first few moments after the reaction is spanned by highly excited longitudinal strings. These strings will later materialize into all the secondaries and the total
energy carried by them, thus leading to a more or less complete distribution of the energy between target and projectile rapidity. In this sense the flux tubes should be interpreted as highly energetic excitations of the QCD vacuum. One might consider such regions as a precursor of a QGP and deconfinement. 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. 6 the position of nucleonic strings (highly energetic strings excited by a target and projectile nucleon) all being produced in a cms-time interval of \( \approx 0.5 \text{ fm/c} \) 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 [25]. In total about 58 strings are produced in the transverse area where the sulphur ion faces the uranium nucleus. This area is about 40 fm\(^2\). If one assumes that the strings are homogeneously distributed (which they are not) this corresponds to a mean area of 0.7 fm\(^2\)/string. Hence all strings are closely packed if the radius of a string is about 0.5 fm (in the Friedberg-Lee model it turns out to be nearly 0.8 fm). 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, a large fraction of the produced \( J/\psi \)'s have to pass it and thus can be affected by this highly excited environment. This 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. If they will get dissociated (in part or all of them) this would lead to a minor or stronger suppression of \( J/\psi \) to be finally observed. Such a conclusion can only be further tested and strengthened if one incorporates these ideas within one of the present microscopic transport algorithms.

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FIG. 1. The average radius $\langle r^2 \rangle^{1/2}$ and the overall mass $M$ of the $c\bar{c}$-state and the $c\bar{q}$-state are shown as a function of the bare charm quark mass $m_q$. 
FIG. 2. The radial distribution of the soliton field and the quark density of the ‘heavy’ quark of the $c\bar{c}$-state is depicted at the initial time (when the constituent quarks are light) and at a very late time (when the masses of the quarks have already increased to their final value).
FIG. 3. The temporal evolution of the total electric field energy inside the string is shown. Up to times $t \approx 13 \text{ fm/c}$ the size of the string is linearly increased to a final length of $\approx 8 \text{ fm}$ and held constant afterwards. At that time $t = 13 \text{ fm/c}$ the $c\bar{c}$-state is inserted into the central region of the string. A nearly linear decrease in time is observed.
FIG. 4. The quark distribution (left side) and the scalar field distribution (right side) at various times in the evolution ($t = 12.8, 13.6, 16.2, 17.4$ fm/c) are shown in a contourplot.
FIG. 5. The quark distribution (left side) and the scalar field distribution (right side) at various times in the evolution ($t = 18.6, 23.2, 26.8, 40.0$ fm/c) are shown in a contourplot.
FIG. 6. The position of nucleonic strings is plotted, all being produced in a cms-time intervall of \( \approx 0.5 \text{ fm/c} \) 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 algorithm [25]. In addition, each string is accompanied by a circle of radius 0.5 fm.