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Mass spectrum of elementary particles in a temperature-dependent model

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Abstract. We show that the temperature-generalization of a popular model of quark-confinement seems to provide a rather interesting insight into the origin of mass of elementary particles: as the universe cooled, there was an era when particles did not have an identity since their masses were variable; the temperature at which the conversion of these 'nomadic' particles into 'elementary' particles took place seems to have been governed by the value of a dimensionless coupling constant C_c. For $C_c = 0{\cdot}001\,(0{\cdot}1)$ this temperature is of the order of 10^9 K (10^{11} K), below which the particle masses do not change.

Keywords. Origin of mass of elementary particles; quark-confinement; temperature-generalized Bethe-Salpeter equation; harmonic oscillator kernel.

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1. Introduction

Mass of an elementary particle is generally taken to be one of its intrinsic attributes, to be experimentally determined. This implies that a physical theory must at the outset have as an input, besides other inputs, as many empirical parameters (masses) as there are elementary particles. Since it is an aim of physics to try and explain natural phenomena on the basis of as few inputs as possible, it is no surprise that physicists have wondered, from time to time, if the large set of observed masses could be obtained from a much smaller set of empirical parameters and a physical principle. Stated differently, these considerations imply the questions: What is the origin of mass of elementary particles? Is there a theory which explains why these particles have the masses that they possess?

Possibly, Einstein's [1] perception of elementary particles as "condensations of the electromagnetic field" marks the conceptual beginning in the above direction. It is interesting to note that Hoyle's [2] concern with the large numbers problem (why does the ratio of the electrical force and the gravitational force between protons equal the ratio of the cosmological distance and the atomic distance?) – spanning a period of 50 years – led him to conclude that the problem is related to the nature of mass. Specifically, demanding the invariance of the line element under a conformal transformation, he asserted that the phenomenon of mass arises from a field interaction; It cannot be an intrinsic property of the particle alone. Essentially similar conclusion was reached by Castell [3] from a different route. In his study of the Bose-Einstein condensation of photons, Castell showed the existence of a second phase for photons which he identified with the matter in the universe. He was thus able to compute a critical temperature (T_c) of $4{\cdot}5 \times 10^{11}$ K, corresponding to a critical mass (m_c) of 39 MeV. Since m_c is about 1/3 of the mass of π meson, Castell interpreted this as providing a connection between the structure of the universe and the properties of elementary particles. One is thus led back to Einstein's perception of massive elementary particles as droplets of photon vapour – with cosmology playing an essential role in the scheme. Indeed, these attempts bring an impressive sweep of ideas to bear on the problem. Nonetheless, it seems to us, we must follow another approach for the most tangible realization of the aim stated above, viz., obtaining the mass spectrum of elementary particles with minimal empirical inputs. The inputs in this approach are the postulate of existence of a few quarks and a mechanism for their confinement. If the quark masses are taken as given, one may vary the parameters of the confining mechanism and attempt to identify the observed particles with the various bound states of these quarks.

Three distinct stages have marked the development of the quark-confinement picture so far. In the first stage, a variety of potentials were investigated with a view to compatibility with the phenomenological features (non-observability of free quarks, rising Regge trajectories, power-law fall off of form factors, etc.) expected of a dynamical model of hadrons [4, 5]. In the next stage, the need to incorporate these features into the framework of relativistic quantum field theory naturally led to the study of the Bethe-Salpeter (BS) equation with a variety of confining kernels [5, 6, 7]. Finally, efforts have recently been directed at investigating the deconfining effects of increasing baryon density [8, 9] and temperature [10, 11, 12].

It is the purpose of this note to present an approach which, in a sense, integrates the three stages of development discussed above. The essence of the approach is to incorporate temperature and density in the BS equation itself, and to study the resulting temperature-generalized equation. We note that the desired generalization can be achieved irrespective of the choice of the confining kernel so long as its use is restricted to the instantaneous (or static) approximation, which indeed is the case in respect of almost all such kernels studied so far [5, 6, 7, 13]. In the next section, therefore, we deal with the temperature-generalization of the BS equation for the bound states of two quarks (treated as spinless) without specifying the kernel used. Our choice of the kernel is dealt with in § 3. The solutions of the resulting equation and their general features are presented in § 4. The final section is devoted to a discussion of our results.

2. Temperature-generalization of the BS equation

The BS equation of interest is:

$$[(P/2+p)^2+m^2-i\varepsilon][(P/2-p)^2+m^2-i\varepsilon]\phi(p) = (-ig^2/\pi^2)\int d^4q\, V(p-q)\phi(q), \qquad (1)$$

which describes the bound states of two equal mass (m) particles. P and p are, respectively, the centre-of-mass and the relative 4-momenta of these particles; g is the strength of interaction between the exchanged particle and each of the constituents of the bound state; the metric used is space-preferred. Equation (1) has been written down in the ladder approximation.

In the instantaneous approximation (which neglects the time of propagation of the exchanged particle) one may set

$$V(p-q)=(2\pi)^{-1}\,V(\mathbf{p}-\mathbf{q}) \tag{2}$$

whence, in the frame $(P=0,E)$, (1) may be written as

$$D(p)\phi(p)=(-ig)^2/(2\pi^3)\int d^4q\,V(\mathbf{p}-\mathbf{q})\phi(q), \tag{3}$$

where

$$D(p)=[(p_0+E/2)^2-(\mathbf{p}^2+m^2)+i\varepsilon]$$
$$\times[(p_0-E/2)^2-(\mathbf{p}^2+m^2)+i\varepsilon]. \tag{4}$$

The choice of $V(\mathbf{p}-\mathbf{q})$ will be spelled out shortly; it does not affect the temperature-generalization of the model, which is carried out as described below.

Since the RHS of (3) is a function of $\mathbf{p}$ alone and not of p_0, we may define

$$D(p)\phi(p)=S(\mathbf{p}), \tag{5}$$

which satisfies the equation

$$S(\mathbf{p})=(-ig^2)/(2\pi^3)\int \mathrm{d}^3q V(\mathbf{p}-\mathbf{q})S(\mathbf{q})I(\mathbf{q}^2), \tag{6}$$

where

$$I(\mathbf{q}^2)=\int \mathrm{d}q_0/D(q). \tag{7}$$

The desired generalization is achieved by affecting the following substitutions [14, 15, 16] in (7)

$$\int \mathrm{d}q_0\to(2\pi i/\beta)\sum_n,\; q_0\to 2\pi ni/\beta \tag{8}$$

where $\beta=(k_\beta T)^{-1}$, k_β being the Boltzmann constant and T the temperature. Note that it is the adoption of the instantaneous approximation which enables us here to temperature-generalize the integral equation under consideration. The use of the recipe (8) in (7) requires some care. It will be seen that the recipe comprises two distinct operations, namely Wick rotation and discretization of the internal energy variable. In affecting the former, one must ensure that the poles and the cuts of the BS equation are not crossed. Thus, for $E>2m$, (7), (4) and (8) yield

$$[I(\mathbf{q}^2)\to\sum=\frac{i\pi A(\mathbf{q}^2)}{\beta E(\mathbf{q}^2+m^2)^{1/2}(\mathbf{q}^2+m^2-E^2/4)}, \tag{9}$$

where

$$A(\mathbf{q}^2)=A_1\coth(A_2)+A_2\coth(|A_1|),$$
$$A_1=(\beta/2)(E/2-\sqrt{\mathbf{q}^2+m^2}), \tag{10}$$

and

$$A_2=(\beta/2)(E/2+\sqrt{\mathbf{q}^2+m^2}).$$

In the limit $T=0$, (9) and (10) yield the standard result of integration over q_0 in

(7), viz.,

$$\int \mathrm{d}p_0/D(q) = i\pi/[2(\mathbf{q}^2+m^2)^{1/2}(\mathbf{q}^2+m^2-E^2/4)]. \tag{11}$$

We note that the case $E < 2m$ has been dealt with in [17].

3. The choice of the kernel

We now deal with the choice of $V(\mathbf{p}-\mathbf{q})$. Among the kernels that have been used in the BS equation to study the problem of confinement of quarks, the one that is interpretable as a harmonic oscillator in the nonrelativistic limit has been by far the most popular [5, 6, 7, 13]. This is easy to understand. Such a kernel not only confines the quarks in a steeply rising potential, yields rising Regge trajectories etc., but also possesses the virtue of being mathematically tractable. Additionally, of course, it enables one to study the oscillator on a quantum scale, an interesting problem in itself. In the following we employ a kernel originally given in [7], but in the form studied in [13]. As has been discussed [7, 13], the static limit of the usual kernel representing the exchange of a massive scalar particle of mass μ yields merely the screened Coulomb or Yukuwa interaction of the form $\exp(-\mu r)/r$. However, higher derivatives with respect to μ of such a kernel soften the singularity at the origin and produce interactions rising asymptotically in configuration space and hence confining in character. Thus, following these authors, we have

$$\begin{aligned} V(\mathbf{p}-\mathbf{q}) &= \lim_{\mu\to 0} d^3/d\mu^3[1/\{(p-q)^2+\mu^2-i\varepsilon\}] \\ &= \lim_{\mu\to 0} 24\mu[\mathbf{p}-\mathbf{q})^2-\mu^2]/[(\mathbf{p}-\mathbf{q})^2+\mu^2]^4 \\ &= 2\pi^2\Delta_p\delta^3(\mathbf{p}-\mathbf{q}). \end{aligned} \tag{12}$$

From (6), (9) and (12) we obtain

$$\Delta_p\chi(\mathbf{p}) = \frac{(\beta E/g^2)(\mathbf{p}^2+m^2)^{1/2}(\mathbf{p}^2+m^2-E^2/4)\chi(\mathbf{p})}{A(\mathbf{p}^2)}, \tag{13}$$

where

$$\chi(\mathbf{p}) = \frac{A(p^2)S(\mathbf{p})}{(p^2+m^2)^{1/2}(p^2+m^2-E^2/4)}. \tag{14}$$

In interpreting (13), the many body origin of the BS equation from which it is derived must be borne in mind. The building blocks of the latter equation are the finite-temperature Green's functions. Recall that the one-particle temperature Green's function describes the motion of one particle added to the many-body system. Since such a function is a propagator, it contains detailed dynamic information; besides, it contains all statistical mechanical information because it is an expectation value in the grand canonical ensemble [18]. Thus, the $T \neq 0$ BS equation may be said to provide a description of the many-body system whose constituents are in a paired state [19]. It will be seen that the chemical potential of the species being considered has not been incorporated into our equation; this is easily achieved by replacing $p^2/(2m)$ by $p^2/(2m) + \tilde{\mu}$ in (9) and (10), where $\tilde{\mu}$ is the chemical potential. In the $\tilde{\mu} = 0$ limit being considered by us, one may regard the $T = 0$ BS equation as describing the bound-state of a 2-particle system in a heated vacuum [20]. Correspondingly, (13) may be said to describe the bound-state system interacting via a temperature-modified potential [21].

4. The solutions

We now proceed to solve (13). Setting

$$\chi(\mathbf{p}) = \chi_l(p)\, Y_{l,m}(\Theta, \Phi)/p, \quad (p = |\mathbf{p}|) \tag{15}$$

the radial part of (13) may be separated as

$$\chi_l''(k) - [l(l+1)/k^2]\chi_l(k) = \frac{(\beta E m^5/g^2)(1+k^2)^{1/2}[k^2 + 1 - E^2/(4m^2)]\chi_l(k)}{A(k^2)}, \tag{16}$$

where k is the dimensionless variable

$$k = p/m. \tag{17}$$

The definitions

$$\begin{aligned} &\alpha = \beta E m^5/g^2 \\ &s_1 = (m\beta/2)[E/(2m) - 1],\ T_1 = \tanh(s_1) \\ &s_2 = (m\beta/2)[E/(2m) + 1],\ T_2 = \tanh(s_2) \\ &\tilde{k}^2 = \beta p^2/(4m) = m\beta k^2/4,\ T_3 = \tanh(\tilde{k}^2) \\ &\gamma^2 = E^2/(4m^2) - 1 > 0 \end{aligned} \tag{18}$$

enable us to write the multiplier of $\chi_l(k)$ on the RHS of (16) as

$$\frac{\alpha(1+k^2)^{1/2}(k^2-\gamma^2)(T_2+T_3)(T_1-T_3)}{s_1 T_1 + s_2 T_2 - T_3^2(s_1 T_2 + s_2 T_1) + Z_1 T_3 + Z_2 \tilde{k}^2} \tag{19}$$

where

$$\begin{aligned} Z_1 &= T_1 T_2(s_1 - s_2) + s_2 - s_1 \\ Z_2 &= T_1 - T_2 - 2T_1 T_2 T_3 + 2T_3 + T_3^2(T_2 - T_1). \end{aligned}$$

From (19) it is transparent that we shall get temperature-independent results so long as s_1 and s_2 are large enough to yield $T_1 = T_2 = 1$. For $m = 1{\cdot}5\,\text{GeV}$ and $C_c = g^2/m^{2\cdot5} = 0{\cdot}1$, this is found to be so for T up to about 10^{11} K. For temperatures exceeding this, we resort to the high-temperature approximation, viz.,

$$T_3 = \tanh(\tilde{k}^2) = \tilde{k}^2. \tag{20}$$

On retaining terms to order $\tilde{k}^2$, (19) is given by

$$\frac{\alpha(1+k^2)^{1/2}(k^2-\gamma^2)[T_1 T_2+\tilde{k}^2(T_1-T_2)]}{[s_1 T_1+s_2 T_2+\tilde{k}^2\{(s_2-s_1)(1-T_1 T_2)+T_2-T_1\}]}$$
$$=\frac{\alpha[-T_1 T_2\gamma^2+k^2\{-(1/4)m\beta\gamma^2(T_1-T_2)+(1-\gamma^2/2)T_1 T_2\}]}{u_1[1+(1/4)u_2 m\beta k^2/u_1]} \quad (21)$$

where we have assumed the kinematics to be nonrelativistic ($k \ll 1$), and have set

$$u_1 = s_1 T_1 + s_2 T_2$$
$$u_2 = (s_2 - s_1)(1 - T_1 T_2) + T_2 - T_1. \quad (22)$$

Substitution of (21) into (16) yields a Schrödinger equation with a rational potential. Note that if it should turn out that

$$u = (u_2/u_1)(m\beta k^2/4) \ll 1 \quad (23)$$

for the values of m, g and T of interest, then (21) can be cast into a particularly simple form:

$$\text{RHS of (16)} = -Q^2 + \Lambda^2 k^2, \quad (24)$$

where

$$Q^2 = \alpha T_1 T_2 \gamma^2/u_1 \quad (25)$$

and

$$\Lambda^2 = (\alpha/u_1)[(1/4)m\beta\gamma^2(u_2 T_1 T_2/u_1 + T_2 - T_1) + T_1 T_2(1-\gamma^2/2)]. \quad (26)$$

From (16) and (24) we therefore obtain

$$\chi_l''(k) + [Q^2 - \Lambda^2 k^2 - l(l+1)/k^2]\chi_l(k) = 0, \quad (27)$$

which is readily recognized as the Schrödinger equation for the (energy-dependent) spherical oscillator potential. The energy eigenvalues of (27) may be numerically obtained from

$$Q^2/(2\Lambda) = (n + 3/2); \quad (28)$$

the eigenfunctions of the equation being given by

$$\chi_{n,l} \sim {}_1F_1(-n_r, l+3/2; \Lambda k^2), \quad (29)$$

where $n_r = (n-1)/2$.

Equation (16)—together with (19) and (27)—enables us to study the model under consideration over a wide range of temperatures. We have done so for several choices of m (e.g., $m = 5, 10$ MeV and 1·5 GeV) and the dimensionless coupling constant $C_c = g/m^{2\cdot 5}$ (e.g., $C_0 = 0\cdot 001$ to 1·0). The consistency of our approximations has been checked by calculating the expectation value of

$$\langle k^2\rangle_n = \langle \beta p^2/(4m)\rangle_n = (2n+3)(\beta m/8\Lambda). \quad (30)$$

From (30) the validity of (20) can be directly checked; also we can now calculate u and hence check the accuracy of the approximation made in (23). Indeed, for a fixed m, one may not expect our approximations to be valid for all values of C_c. The general features of the temperature-generalized model considered by us are:

(a) For any chosen value of m, there exists an upper limit for C_c up to which our approximations, viz., (20) and (23) are valid.
(b) For any admissible values of m and C_c there exists a certain temperature T up to which the model admits solutions which differ little from the corresponding solutions for the $T = 0$ model. For $m = 1{\cdot}5$ GeV and $C_c = 0{\cdot}001$, for which we present our results in table 1 as an illustration, T_c is of the order of 10^9 K.
(c) For $T > T_c$, the energy eigenvalues of the system increase with temperature. This result is consistent with the accepted view that the force between the quarks, or the "spring constant", increases with their separation since—after a threshold temperature (T_c)—one would expect the separation between the constituents of our system to increase.
(d) We have restricted our study of the model up to a temperature T for which $k_B T < m$, since it would be clearly inadmissible to trust the results of a nonrelativistic approach beyond this temperature.
(e) For values of C_c for which the approximation made in (23) does not hold, one must indeed resort to a study of the Schrödinger equation with the rational potential, viz., (21).

4. Concluding remarks

a) That the use of the recipe given in (8) enables one to obtain the analytic expression for any $T \neq 0$ Feynman diagram from its $T = 0$ counterpart is familiar to most workers

Table 1. The first ten energy eigenvalues of the temperature-generalized model for $m = 1500$ MeV and $C_c = 0{\cdot}001$ at different temperatures. The final entries in each column correspond, for the $n = 0$ state, to the values of $\langle \tilde{k}^2 \rangle$ and u [see eqs (20) and (23)], respectively.

Temp(K)	$0 - 1 \times 10^9$	1×10^{10}	1×10^{11}	1×10^{12}	1×10^{13}
n			E_n(MeV)		
0	3003·179	3004·164	3008·858	3019·051	3040·895
1	3005·294	3006·012	3012·445	3026·746	3057·331
2	3007·406	3007·837	3015·569	3033·434	3071·579
3	3009·516	3009·737	3018·404	3039·492	3084·455
4	3011·622	3011·723	3021·036	3045·103	3096·356
5	3013·726	3013·767	3023·516	3050·373	3107·510
6	3015·826	3015·842	3025·874	3055·369	3118·062
7	3017·924	3017·930	3028·136	3060·139	3128·131
8	3020·018	3020·020	3030·316	3064·719	3137·775
9	3022·110	3022·111	3032·429	3069·133	3147·058
$\langle \tilde{k}^2 \rangle$,		0·42,	0·07,	0·01,	0·003,
u		0·07	0·05	0·01	0·005,

in condensed matter physics/finite temperature field theory (for a brief description of the origin of the recipe, we refer the reader to [22]). What we have been concerned with in this paper is the temperature-generalization of an *integral equation*, which incorporates not only the sum of a class of Feynman diagrams, but also the (to be determined) Bethe-Salpeter wavefunction. That the recipe should suffice for this purpose is, perhaps, not clear. One might therefore desire that an *ab initio* derivation of the $T \neq 0$ Bethe-Salpeter equation be attempted to check if, in fact, the use of the recipe yields the same equation. Such an exercise has already been done for the Bethe-Salpeter equation for the bound states of a many electron-proton system [19]: it relies on the familiar property of periodicity of the Green's functions in the variable $\tau(= i\beta)$, and the homogeneity of the medium; it does not require any additional assumptions. The result is, of course, that the $T \neq 0$ equation derived from first principles agrees with the $T \neq 0$ equation obtained via the well-known recipe used by us [22]. We note that a careful interpretation of the thermal Bethe-Salpeter equation has been given in the text.

b) It seems worthwhile to draw attention to somewhat different frameworks within which the effect of temperature on hadron properties has been recently investigated. Bernard and Meissner [12], for example, use the topological chiral soliton model based on a nonlinear meson theory of interacting pions (ρ-and w- mesons) for the description of the nucleon. The nucleon properties are determined by the mesonic parameters. For the temperature-dependence of these meson parameters, they use rigorous results from chiral perturbanon theory as well as predictions from the Nambu-Jona-Lasinio model. Another approach to the problem employs the Monte Carlo methods to study SU(2) Yang-Mills theory on a lattice at finite temperature [25, 26, 27]. In each of these investigations a physical property of primary interest is the temperature which marks the transition between the unconfined and the confined phases. Indeed, the signal of this transition is different in the two approaches. Thus, Bernard and Meissner [12] find that close to a certain temperature, the effective nucleon mass decreases sharply and the baryon charge is spread out. Eventually, nucleons lose their individual character, which is taken to be the signal of deconfinement. Kuti *et al* [26], on the other hand, distinguish the two phases by the different behaviours of a correlation function defined in terms of thermal Wilson loops. Interestingly enough, the transition temperatures in the two approaches are, respectively, 180 MeV and 160 ± 30 MeV. These translate to about 10^{12} K. If we interpret the temperature regime in which our bound state system has non-constant energy eigenvalues as the unconfined phase, the corresponding temperature in our approach is of the order of 10^{11} K for $C_c = 0{\cdot}1$.

c) The incorporation of chemical potential into the model has been discussed in the text; it is easy to see that while it will shift the spectrum of bound state eigenvalues, it will not alter the general features of the solutions obtained by us. Work is currently in progress to study the effect of temperature and chemical potential in bringing about the deconfining phase.

d) That the elementary particles had a 'nomadic' phase before they found their niche seems to be a rather interesting result; see also [12] in this connection. We note that we have been led to it by following an approach which incorporates temperature in the dynamics itself. It seems pertinent to conclude by drawing attention to what seem to be significant results obtained by following the same approach in another field, viz., astrophysical plasmas [21–24].

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