Properties of vector mesons at finite temperature: Effective Lagrangian approach

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We examine the properties of ρ mesons at finite temperature T with the use of an effective chiral Lagrangian. The effective Lagrangian includes vector and axial-vector mesons as massive Yang-Mills fields of the chiral symmetry. It is shown that at the leading order of the temperature the effective mass of the ρ meson is not changed and only a mixing in vector and axial-vector correlators takes place.

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I. INTRODUCTION

It is expected that at very high temperatures and/or densities hadronic matter undergoes a phase transition or crossover into a plasma phase composed of weakly interacting quarks and gluons [1]. It may be possible to produce and observe this new phase of hadronic matter in relativistic collisions between two heavy nuclei. Experiments are now being carried out at the BNL Alternate Gradiant Synchrotron (AGS) and at the CERN Super Proton Synchrotron (SPS), and in the future there will be experiments for which one anticipates that very high energy-densities will be achieved; they will be high enough to form a quark-gluon plasma. In the plasma phase of hadronic matter the spontaneously broken chiral symmetry is also expected to be restored. The restoration of the chiral symmetry can be characterized by the vanishing of the quark condensate, which is known as the order parameter of the phase transition. It has been shown by the numerical simulation of QCD formulated on the lattice [2] and also from some model calculations [3] that the quark condensate diminishes with rising temperature.

We are interested in the phenomena which arise prior to the chiral transition in hot hadronic matter. The hadronic phase below the phase transition temperature can be described by hadronic degrees of freedom: pions, kaons, vector and axial-vector mesons, etc. The precursor phenomena of the chiral phase transition then can be observed by scrutinizing those properties of hadrons which would be affected by the change of the condensate at finite temperature. It is therefore desirable to have a direct connection between the properties of hadrons and the quark condensate in order to study the chiral phase transition in hot hadronic matter. Of particular interest are the properties of vector mesons at finite temperature since model calculations show a definite relation between chiral symmetry restoration at finite temperature and the mass of the vector meson [4]. Moreover, it is possible to observe the change of the mass in the dilepton spectrum produced from hot hadronic matter, since the peak in the spectrum corresponding to the vector meson will be shifted.

The properties of vector mesons and the dilepton spectrum at finite temperature have been studied in various ways [4–12]. Recently, it was realized by Dey *et. al.* [6] that at the lowest order of T^2/f_π^2, where the pion decay constant $f_\pi = 93$ MeV, the mass of the vector meson is not changed but only mixing between vector and axial-vector correlators takes place. Since the isospin mixing at finite temperature is obtained from PCAC (partial conservation of axial-vector current) and current algebra, it has to be satisfied in the low temperature limit of any model calculations as long as the same symmetry properties are preserved [7,8,12].

In this paper we show that the isospin mixing phenomenon at finite temperature is obtained from an effective Lagrangian approach. In Sec. II we introduce an effective chiral Lagrangian which includes vector and axial-vector mesons as massive Yang-Mills fields of the chiral group and photon fields via the vector meson dominance assumption. Parameters in the effective Lagrangian are chosen to reproduce the Kawarabayashi-Suzuki-Riazuddin-Fayyadzuddin (KSRF) relation, universality of vector meson coupling and Weinberg's mass relation. The properties of the ρ meson at finite temperature are studied with the effective Lagrangian. It is shown that, at the T^2 order, there is a mixing in vector and axial-vector correlators and also there is no change in the effective mass of the vector meson. In Sec. III we introduce an extra nonminimal coupling term to reproduce the experimental values for the mass and width of the axial-vector meson. With the new term we cannot get the exact mixing effect in the vector and axial-vector channels at finite temperature and have a thermal correction to the mass proportional to T^2. However, the effective mass of the ρ meson is not changed much at the leading order of the temperature.

II. ISOSPIN MIXING AT FINITE TEMPERATURE

We consider an effective chiral Lagrangian with vector and axial-vector meson fields which are introduced as massive Yang-Mills fields [13]:

$$\mathscr{L} = \tfrac{1}{4} f_\pi^2 \mathrm{Tr}[D_\mu U D^\mu U^\dagger] - \tfrac{1}{2}\, \mathrm{Tr}[F^L_{\mu\nu} F^{L\mu\nu} + F^R_{\mu\nu} F^{R\mu\nu}] + m_0^2 \mathrm{Tr}[A^L_\mu A^{L\mu} + A^R_\mu A^{R\mu}] - i\xi \mathrm{Tr}[D_\mu U D_\nu U^\dagger F^{L\mu\nu} + D_\mu U^\dagger D_\nu U F^{R\mu\nu}], \tag{1}$$

where U is related to the pseudoscalar fields ϕ by

$$U = \exp\left[\frac{i\sqrt{2}}{f_\pi}\phi\right], \quad \phi = \sum_{a=1}^{3} \phi_a \frac{\tau_a}{\sqrt{2}}, \tag{2}$$

and $A^L_\mu (A^R_\mu)$ are left- (right-)handed vector fields. The covariant derivative acting on U is given by

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$$D_\mu U = \partial_\mu U - igA^L_\mu U - igUA^R_\mu, \tag{3}$$

and $F^L_{\mu\nu}(F^R_{\mu\nu})$ is the field tensor of left- (right-)handed vector fields. The A^L_μ and A^R_μ can be written in terms of vector (V_μ) and axial-vector fields (A_μ) as

$$A^L_\mu = \tfrac{1}{2}(V_\mu - A_\mu), \quad A^R_\mu = \tfrac{1}{2}(V_\mu + A_\mu). \tag{4}$$

The Lagrangian should be diagonalized by definitions

$$A_\mu \to A_\mu + \frac{gf_\pi}{\sqrt{2}m_0^2}\left(\partial_\mu\phi - i\frac{g}{2}[V_\mu,\phi]\right),$$

$$\phi \to Z^{-1}\phi, \quad f_\pi \to Z^{-1}f_\pi. \tag{5}$$

In terms of new fields we find

$$\mathscr{L}^{(2)} = \tfrac{1}{2}\mathrm{Tr}\left(\partial_\mu\phi - i\frac{g}{2}[V_\mu,\phi]\right)^2 - \tfrac{1}{4}\ \mathrm{Tr}[F^V_{\mu\nu}F^{V\mu\nu} + F^A_{\mu\nu}F^{A\mu\nu}] + \tfrac{1}{2}m_\rho^2\mathrm{Tr}V_\mu^2 + \tfrac{1}{2}m_a^2\mathrm{Tr}A_\mu^2, \tag{6}$$

where we use

$$Z^2 = \left[1 - \frac{g^2f_\pi^2}{2m_\rho^2}\right] = \frac{m_\rho^2}{m_a^2}, \tag{7}$$

and the masses of the vector and axial-vector meson are given by

$$m_V^2 = m_0^2, \quad m_A^2 = m_0^2/Z^2. \tag{8}$$

When we choose $Z^2 = 1/2$ we have the KSRF relation, $m_\rho^2 = g^2f_\pi^2$, and Weinberg's mass relation, $m_a^2 = 2m_\rho^2$.

The decay rate for $\rho \to \pi\pi$ is given by

$$\Gamma(\rho \to \pi\pi) = \frac{1}{6\pi m_\rho^2}|q_\pi|^3 g^2_{\rho\pi\pi}, \tag{9}$$

with

$$g_{\rho\pi\pi} = \frac{g}{4\sqrt{2}}(3 + 2g\xi). \tag{10}$$

To satisfy the universality of vector meson coupling $g_{\rho\pi\pi} = g/\sqrt{2}$, we choose $2g\xi = 1$. (There is a factor $1/\sqrt{2}$ because we use convention $V_\mu = \tau^a/\sqrt{2}\,\rho^a_\mu$, where a is the isospin index.) The coupling constant g can be determined from the width of the ρ meson, $\Gamma^{\mathrm{expt}}(\rho \to \pi\pi) = 150$ MeV.

With the effective Lagrangian we study those properties of vector mesons at finite temperature which will be modified due to the interaction with particles in a heat bath. Here we assume that the known hadronic interactions can be used at finite temperature and describe the interactions among particles in hot hadronic matter. To study the modification we calculate the self-energy of the vector meson which is defined by the difference between the inverse of the in-medium propagator $D_{\mu\nu}$ and that of the vacuum propagator $D^0_{\mu\nu}$ [14]:

$$\Pi_{\mu\nu} = D^{-1}_{\mu\nu} - D^{-1}_{0\mu\nu}. \tag{11}$$

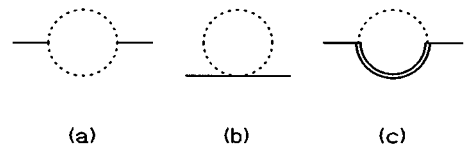

FIG. 1. One loop diagrams for the self-energy of the ρ meson. The dotted, solid, and double solid lines denote, respectively, the pion, ρ meson, and a_1 meson.

Since the self-energy of the vector meson is symmetric and transverse it can be written in terms of the longitudinal ($P^L_{\mu\nu}$) and transversal ($P^T_{\mu\nu}$) projection tensors as

$$\Pi_{\mu\nu} = GP^T_{\mu\nu} + FP^L_{\mu\nu}, \tag{12}$$

where $P^T_{00} = P^T_{0i} = P^T_{i0} = 0$, $P^T_{ij} = \delta_{ij} - k_ik_j/k^2$, and $P^L_{\mu\nu} = K^\mu K^\nu/K^2 - g_{\mu\nu} - P^T_{\mu\nu}$. In this paper capital letters are used for four-dimensional momenta: $K = (k_0, \vec{k})$, $k = |\vec{k}|$, and $K^2 = k_0^2 - \vec{k}^2$. The functions F and G are given by

$$F(k_0,\vec{k}) = \frac{K^2}{k^2}\Pi_{00}(k_0,\vec{k}),$$

$$G(k_0,\vec{k}) = -\tfrac{1}{2}[\Pi^\mu_\mu(k_0,\vec{k}) + F(k_0,\vec{k})]. \tag{13}$$

The propagator of the vector meson in the medium can be written as

$$\mathscr{D}_{\mu\nu} = -\frac{P_L^{\mu\nu}}{K^2 - m_\rho^2 - F} - \frac{P_T^{\mu\nu}}{K^2 - m_\rho^2 - G} - \frac{K^\mu K^\nu}{m_\rho^2 K^2}. \tag{14}$$

We consider the corrections that come from the interaction with thermal pions only. Particles with mass M generate contributions of order $\exp(-M/T)$ which are exponentially suppressed compared with the effect from thermal pions. The interaction with pions generates corrections in a power series controlled by the expansion parameter $\sim T^2/f_\pi^2$. The thermal corrections can be obtained in a systematic way. In this respect, it is similar to the calculation of the loop corrections in chiral perturbation theory [15]. The self-energy of the vector meson in hot hadronic matter then can be expanded in powers of T^2/f_π^2 as

$$\Pi_{\mu\nu}(k_0,k;T) = \Pi^{(1)}_{\mu\nu}(k_0,k;T^2/f_\pi^2) + \Pi^{(2)}_{\mu\nu}(k_0,k;T^4/f_\pi^4) + \cdots. \tag{15}$$

The leading contributions can be obtained from one-loop diagrams in Fig. 1. Even in the presence of the a_1 meson in a loop we can still make a systematic expansion. At the leading order of the temperature the contribution from diagram Fig. 1(c) can be written as

$$\Pi_{\mu\nu}^{(c)}(k_0,k)=\frac{1}{2f_\pi^2}K^2(K^2g_{\mu\nu}-K_\mu K_\nu)T \times\sum\int\frac{d^3p}{(2\pi)^3}\frac{1}{P^2}\frac{1}{(P-K)^2-m_a^2}. \quad (16)$$

The full expression is given in the Appendix. For the axial-vector meson in the loop with momentum Q, we have $1/(Q^2-m_a^2)$ in which the momentum Q is given by the momentum of the thermal pion P, and the external vector meson K. When we sum over the Matsubara frequencies of thermal pions we get

$$T\sum\int\frac{d^3p}{(2\pi)^3}\frac{1}{P^2}\frac{1}{(P-K)^2-m_a^2} \sim\frac{1}{K^2-m_a^2}\int\frac{p\ dp}{e^{p/T}-1}\left(1+\frac{p^2}{K^2-m_a^2}+\frac{p^4}{(K^2-m_a^2)^2}+\cdots\right), \quad (17)$$

where we assume that

$$\frac{p^2}{K^2-m_a^2}\sim\frac{p^2}{m_\rho^2}\ll 1. \quad (18)$$

When we use the KSRF relation $m_\rho^2=g^2f_\pi^2$ we can still expand the correction from the $\pi-a_1$ loop in powers of T^2/f_π^2.

The self-energy includes both real and imaginary parts. While the real part describes the excitation spectrum in hot matter, the imaginary part is related to the dissipative properties of the system. The imaginary part is obtained when we analytically continue the Matsubara frequency $k_0\to\omega+i\epsilon$, $\epsilon\to 0^+$ and use

$$\frac{1}{x-x_0\pm i\epsilon}=\mathscr{P}\frac{1}{x-x_0}\mp i\pi\delta(x-x_0), \quad (19)$$

where $\mathscr{P}$ is the principal value. For rho mesons at finite temperature the imaginary terms appear in diagrams (a) and (c) only when the momentum of the thermal pion p lies in the interval

$$\tfrac{1}{2}|\omega-k|\leqslant p\leqslant\tfrac{1}{2}(\omega+k), \quad (20)$$

where we use the relation $m_a^2=2m_\rho^2$ and ω and k are energy and momentum of the ρ meson, respectively. Since $|\omega-k|\geqslant m_\rho$ this condition is not satisfied as long as the soft momentum limit for thermal pions is taken as (18). Thus there is no temperature-dependent correction to the imaginary part of the self-energy in the limit we take.

For the real part we get, at the leading order of T^2/f_π^2,

$$\mathrm{Re}[F^\pi(k_0,k)]=\frac{K^2}{k^2}g^2f_\pi^2\left[\frac{k_0^2}{K^2}-1-\frac{k^2}{2m_\rho^2}\right]\frac{T^2}{12f_\pi^2},$$

$$\mathrm{Re}[F^{a_1}(k_0,k)]=K^2\left(1+\frac{m_a^2}{K^2-m_a^2}\right)\frac{T^2}{24f_\pi^2}, \quad (21)$$

where F^π is the contribution comes from the pion loops in Figs. 1(a) and 1(b) and F^{a_1} is from the $\pi-a_1$ loop in Fig. 1(c). The first two terms in F^π have been obtained from a calculation made only with charged pions and neutral vector mesons [9]. Including a_1 mesons we have a contribution from the $\pi-a_1$ loop and the last term in F^π. When we use the relations

$$g^2f_\pi^2=m_\rho^2, \quad m_a^2=2m_\rho^2, \quad (22)$$

there is an exact cancellation of the last term in F^π and the first term in F^{a_1}, and finally we have

$$\mathrm{Re}[F(k_0,k)]=\mathrm{Re}[F^\pi(k_0,k)]+\mathrm{Re}[F^{a_1}(k_0,k)] =\left(m_\rho^2+\frac{m_\rho^4}{K^2-m_a^2}\right)\frac{T^2}{6f_\pi^2}. \quad (23)$$

Inserting (23) into (14) we can show that there is an isospin-mixing effect at the leading order of the temperature in the longitudinal mode of the vector meson propagator:

$$\frac{1}{K^2-m_\rho^2-F}=\frac{1}{K^2-m_\rho^2}+\frac{1}{K^2-m_\rho^2}F\frac{1}{K^2-m_\rho^2}+\cdots =\left(1-\frac{T^2}{6f_\pi^2}\right)\frac{1}{K^2-m_\rho^2}+\frac{T^2}{6f_\pi^2}\frac{1}{K^2-m_a^2} +O(T^4/f_\pi^4). \quad (24)$$

For the transverse component we have

$$\mathrm{Re}[G(k_0,k)]=-\frac{1}{2}\left[g^2f_\pi^2\frac{T^2}{12f_\pi^2}-g^2f_\pi^2\left(4-\frac{3K^2}{2m_\rho^2}\right)\frac{T^2}{12f_\pi^2} -3K^2\left(1+\frac{m_a^2}{K^2-m_a^2}\right)\frac{T^2}{24f_\pi^2}+\mathrm{Re}[F(k_0,k)]\right] =\mathrm{Re}[F(k_0,k)]. \quad (25)$$

Thus the transverse and longitudinal components are the same at the T^2 order and the mixing effect appears in both modes of the vector meson propagator at finite temperature.

III. EFFECTIVE MASS OF THE VECTOR MESON

Even though the effective Lagrangian we used satisfies the universality and KSRF relations, the Lagrangian could not reproduce the experimental values for the mass and decay width of the a_1 meson. For given parameters with $2g\xi=1$ and $Z^2=1/2$, we have $m_a=\sqrt{2}\ m_\rho=1089$ MeV and $\Gamma(a_1\to\pi\rho)=53$ MeV, while experiments show that $m_a^{\mathrm{expt}}=1260$ MeV and $\Gamma^{\mathrm{expt}}(a_1\to\pi\rho)=400$ MeV. It is possible to describe well the masses and widths of the vector and axial-vector mesons by adding an extra nonminimal coupling term [16]

$$\mathscr{L}_\sigma=\sigma\ \mathrm{Tr}[F_{\mu\nu}^LUF^{R\mu\nu}U^\dagger], \quad (26)$$

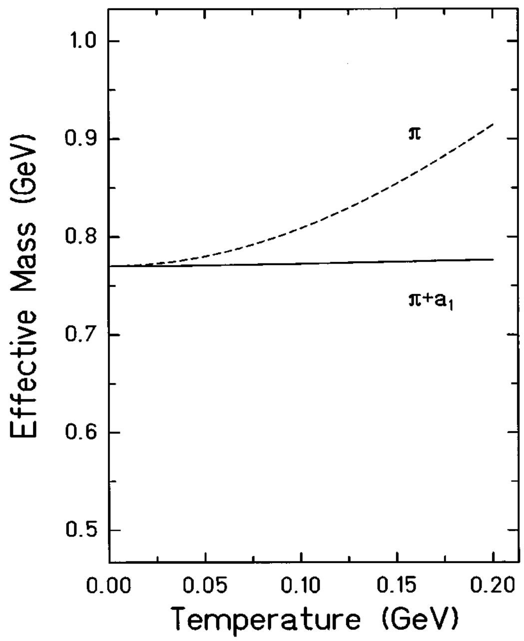

FIG. 2. Effective mass of the ρ meson at finite temperature. The dashed line is the result from the calculation with pions and rho mesons. The solid line is the result obtained when a_1 mesons are included.

with parameters $g=10.3063, \sigma=0.3405, \xi=0.4473$ [10]. However, the Lagrangian does not satisfy the universality, $g_{\rho\pi\pi}=3.06(g/\sqrt{2})$ and KSRF relations, $m_\rho^2=1.63g^2f_\pi^2$.

With these parameters we have

$$\mathrm{Re}[F(k_0,k;T)]=g^2f_\pi^2\frac{T^2}{12f_\pi^2}+K^2\left(2\eta_2^2-\frac{\lambda g^2f_\pi^2}{m_\rho^2}\right)\frac{T^2}{12f_\pi^2}+\eta_2^2K^2\frac{m_a^2}{K^2-m_a^2}\frac{T^2}{6f_\pi^2}, \tag{27}$$

where η_2 and λ are given in the Appendix. By introducing an extra term, the mixing bewteen the vector and axial-vector correlators is not so exact as that shown from the calculation based on current algebra and PCAC, and the effective mass of the ρ meson has a T^2 dependent correction.

We calculate the effective mass of the vector meson which is defined as the pole position of the propagator with zero three-momentum. Since there is no distinction between the longitudinal and transverse modes in the limit $\vec{k}\to 0$, the effective mass of the vector meson is obtained from the equation

$$k_0^2-m_\rho^2-\mathrm{Re}[F(k_0,\vec{k}\to 0;T)]=0. \tag{28}$$

From (27) and (28) we see that the effective mass of the ρ meson is not changed with temperature when $\sigma=0$, $2g\xi=1$, and $Z^2=1/2$. For $\sigma\neq 0$, we have a correction proportional to T^2 that is shown in Fig. 2. The dashed line is the result obtained from the calculation only with pions and rho mesons, and the solid line is that obtained by including a_1 mesons in the effective Lagrangian. We find that the effective mass of the vector meson is not changed much with temperature at the T^2 order [17] and the increase due to pion loops is almost canceled out when we include a_1 mesons.

IV. CONCLUSION

We studied the properties of the vector meson at finite temperature with an effective chiral Lagrangian in which vector and axial-vector mesons are introduced as massive Yang-Mills fields. It is shown that, at the leading order of the temperature, isospin mixing in vector and axial-vector correlators takes place and the effective mass of the vector meson is not changed. When we include an extra term $\mathscr{L}_\sigma$ in the Lagrangian to fit the experimental values for the mass and width of the a_1 meson, the isospin mixing does not have the same form as that shown from model independent calculation, and there is a T^2 dependent correction to the effective mass of the ρ meson. However, the effect turns out to be very small. The increase of the effective mass due to pion-loop corrections is almost canceled by the contribution from the $\pi-a_1$ loop.

This result implies that the change in the effective mass of the vector meson at finite temperature cannot be observed unless the temperature is very close to the critical value for the phase transition. Instead, at low temperature there is an appreciable reduction in the coupling constant of the external vector current to vector mesons because of the mixing effect [12]. The reduction in the coupling constant, which has stronger dependence on the temperature than the shift of the peak position in the spectrum, leads to a suppression in the production rates of photons and dileptons from hot hadronic matter [18].

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APPENDIX: ρ MESON SELF-ENERGY FROM ONE-LOOP DIAGRAMS

The self-energy of the ρ meson is given by

$$\Pi_{\mu\nu}^{(a)}(k_0,k)=-\frac{1}{2}g^2T\sum\int\frac{d^3p}{(2\pi)^3}\frac{(2P_\mu-K_\mu)(2P_\nu-K_\nu)}{P^2(P-K)^2}, \tag{A1}$$

$$\Pi_{\mu\nu}^{(b)}(k_0,k)=\left[g_{\mu\nu}-\frac{\lambda}{m_\rho^2}(K^2g_{\mu\nu}-K_\mu K_\nu)\right]g^2T\times\sum\int\frac{d^3p}{(2\pi)^3}\frac{1}{P^2}, \tag{A2}$$

$$\Pi^{(c)}_{\mu\nu}(k_0,k)=\frac{2}{f_\pi^2}T\sum\int\frac{d^3p}{(2\pi)^3}\frac{1}{P^2}\frac{1}{(P-K)^2-m_a^2}\times\Big\{\eta_2^2K^2(K^2g_{\mu\nu}-K_\mu K_\nu)+2\eta_2\bar{\eta}(K\cdot P)\times(K^2g_{\mu\nu}-K_\mu K_\nu)+\bar{\eta}^2[(K\cdot P)^2g_{\mu\nu}-(K\cdot P)\times(P_\mu K_\nu+P_\nu K_\mu)+K^2P_\mu P_\nu]+\frac{\eta_1^2}{m_a^2}[(K^2)^2P_\mu P_\nu-K^2(K\cdot P)\times(P_\mu K_\nu+P_\nu K_\mu)+(K\cdot P)^2K_\mu K_\nu]\Big\}, \qquad (A3)$$

where the superscript a,b,c denote the contributions from Figs. 1(a)–(c), respectively, and the

$$\eta_1=\frac{g^2f_\pi^2}{2m_\rho^2}\left(\frac{1-\sigma}{1+\sigma}\right)^{1/2}+\frac{2g\xi}{\sqrt{1+\sigma}}\left(\frac{1-\sigma}{1+\sigma}\right)\frac{m_\rho^2}{m_a^2}, \qquad (A4)$$

$$\eta_2=\frac{g^2f_\pi^2}{2m_\rho^2}\left(\frac{1+\sigma}{1-\sigma}\right)^{1/2}-\frac{2\sigma}{\sqrt{1-\sigma^2}}, \qquad (A5)$$

$$\lambda=\frac{g^2f_\pi^2}{2m_\rho^2}\left(\frac{1+\sigma}{1-\sigma}\right)-\frac{4\sigma}{1-\sigma}\left(1-\frac{m_\rho^2}{g^2f_\pi^2}\right), \qquad (A6)$$

and $\bar{\eta}=\eta_1-\eta_2$. When $\sigma=0$, $2g\xi=1$ and $Z^2=1/2$ we have $\eta_2=1/2$ and $\lambda=1/2$.

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