Anomalous properties and quantized massive gauge fields in high-Tc cuprates

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Abstract. Quantized massive gauge fields have been introduced as collective modes, which might contain effects of spin fluctuations, charge fluctuations, and phonons, around doped holes from the viewpoint of adjoining interactions. One has discussed the relationship among anomalous properties in high-Tc cuprates from the standpoint of the effect of quantized massive gauge fields around doped holes and the restoration of spontaneous symmetry breaking. It is suggested strongly that the quantized massive gauge fields are mediating Cooper pairing in high-Tc cuprates.

1. Introduction
For high Tc cuprates, the transition temperature at which they become superconducting is much higher than in other materials. The mechanism that gives rise to high-transition-temperature superconductivity has been debated since the discovery of the phenomenon. In underdoped cuprates, it is known that there is considerable evidence that two aspects of the pseudogaps, small pseudogap (SPG) and large pseudogap (LPG), in the magnetic and electronic excitation spectra in the normal state [1,2]. Anomalous normal-state transport in high-Tc cuprates remains controversial subject. For example, the electrical resistivity is linear in temperature in a wide range of temperature above Tc. Some mechanisms for the anomalous normal state transport property have been proposed [11-13]. In addition, the mechanism of anisotropic normal state transport in high-Tc cuprates remains a controversial subject. Various models, which may explain the anisotropic scattering property of a hole in Cu-O planes, have been proposed [14-18]. A central concept is the presence of hot spots near the (π, 0) and (0, π) points. Around these hot spots, the electron lifetime is unusual short and has an anomalous temperature dependence. In cold spots, which are small regions near the zone diagonal on the Fermi surface, the electron lifetime is much longer than elsewhere on the Fermi surface. This view is consistent with experiments [19]. Recently, Ando et al. [20] have observed that the magnitude of Hall coefficient markedly decreases with increasing temperature in lightly doped LaSrCuO. This result suggests that the effective carrier density becomes larger than the nominal hole density at elevated temperatures, which might be related to the T-evolution of the Fermi arc in underdoped cuprates [3]. Recently, Hwang et al.[21] have measured the so-called optical self-energy of a bismuth-containing copper oxide, and have found that the optical infrared-spectra has a broad and featureless distribution (the broad background), whose intensity decreases gradually as the temperature increases. The broad background is a universal property of the copper-oxygen plane and provides a good candidate signature of the 'glue' that binds the holes into Cooper pairs.
In this paper, we will discuss the relationship among large pseudogap (LGP), the restoration of the spontaneous symmetry breaking, small pseudogap (SPG), and anomalous properties in high-T_c cuprates, from the standpoint of the effect of the quantized massive gauge fields around the holes, which have been introduced theoretically in the path-integral method by the present author [8-10].

2. A model system and quantized massive gauge fields

The SPG, whose corresponding temperature is \( T^* \), is a characteristic energy comparable to the superconducting (SC)-state gap and the LPG is 3 to 4 times larger. The SPG develops progressively below the temperature \( T^* \) and then evolves into the SC-state gap at \( T_c \). This means that the SPG is some type of precursor with superconductivity. On the other hand, the LPG, which remains open even at high temperatures above the mean-field critical temperature \( T_{mc} \), seems to be related to the crossover behavior of magnetic susceptibility at approximately \( T_{max}( > T_{mc}) \), which arises from the gradual development of antiferromagnetic (AF) short-range states. In addition, it has been reported on \( YBCO_{6.6} \) and \( Bi_{2212} \) cuprates that the spin part of the Knight shift decreases linearly upon cooling from \( T_{mK} \) to \( T^* \), and then decreases remarkably upon cooling from \( T^* \) [22]. Taking into account that the symmetry in the undoped \((2 + 1)\)-dimensional quantum antiferromagnet is invariant under local \( SU(2) \) [23], we think that the perturbing gauge fields \( A_{\mu}^a \) introduced by the hole has a local \( SU(2) \) symmetry. That is, since the distortion around the hole is due to strong many-body effects, the distortion effects around the hole is non-linear (Yang-Mill fields-like). Then it is assumed that \( SU(2) \) gauge fields \( A_{\mu}^a \) are spontaneously broken through the Anderson-Higgs mechanism in a way similar to the breaking of the antiferromagnetic symmetry around the hole. We set the symmetry breaking \( \langle 0|\phi_a|0 \rangle = \langle 0, 0, \mu(\hat{k}_F) \rangle \) of the Bose field \( \phi_a \) in the Lagrangian density as follows,

\[
L = \frac{1}{2} \left( \partial_i N_i^j - g_1 \varepsilon_{abc} \varepsilon_{ijk} A_i^a N_j^k \right)^2 + \psi^+ (i\partial_0 - g_2 T_a A_0^a) \psi \\
- \frac{1}{2m} \psi^+ \left( i \nabla - g_2 T_a A_0^a \right) \psi \\
- \frac{1}{4} \left( \partial_\mu A_\mu^a - \partial_\mu A_\mu^a + g_3 \varepsilon_{abc} A_\mu^b A_\mu^c \right)^2 + \frac{1}{2} \left( \partial_\mu \phi_a - g_4 \varepsilon_{abc} A_\mu^b \phi_c \right)^2 \\
- \lambda^2 \left( \phi_a \phi_a - \mu^2 \right)^2 \tag{1}
\]

\( \hat{k}_F \) is the vector of the Fermi momentum. After the symmetry breaking \( \langle 0|\phi_a|0 \rangle = \langle 0, 0, \mu(\hat{k}_F) \rangle \), that is, transition of fields \( \phi_a \) with a zero asymptotics at infinity,

\[
\langle \phi_1, \phi_2, \phi_3 \rangle \longrightarrow \langle \phi_1, \phi_2, \mu(\hat{k}_F) + \phi_3 \rangle
\]

makes the isotopic-symmetry breaking explicit, and we can obtain the effective Lagrangian density, \( \mathcal{L}_{eff} \), [6-10]. That is, \( \langle 0|\phi_3|0 \rangle \) can be regarded as a kind of the disorder parameter [24]. The value, \( \mu(\hat{k}_F) \), of the symmetry breaking depends strongly on the direction of Fermi momentum, \( \hat{k}_F \), on the Fermi surface. Furthermore, the value \( \mu(\hat{k}_F) \), is much correlated to the gap-energy of the high energy pseudogap. If the value of the high-energy pseudogap is related to the strength of the antiferromagnetic short-range order [2], the distortion, which is induced by the doped hole, becomes larger as the hole is doped in the state of the larger gap-energy of the high-energy pseudogap. Since the value, \( \mu(\hat{k}_F) \), means the strength of the distortion induced
by the doped hole, the value, $\mu(\tilde{k}_p)$, is higher around the hot spot.

$$L_{\text{eff}} = \frac{1}{2} \left( \partial_i N^i_c - g_1 \varepsilon_{abc} \varepsilon_{ijk} A^b_i N^k_a \right)^2 + \psi^+ \left( i \partial_0 - g_2 T_a A^0_a \right) \psi - \frac{1}{2m} \psi^+ \left( i \nabla - g_2 T_a A^a_{\mu \neq 0} \right)^2 \psi - \frac{1}{4} \left( \partial_\mu A^a_\mu - \partial_\mu A^a_\mu + g_3 \varepsilon_{abc} A^b_\mu A^c_\nu \right)^2 + \frac{1}{2} \left( \partial_\mu \phi_a - g_4 \varepsilon_{abc} A^b_\mu \phi_c \right)^2 + \frac{1}{2} \mu_1^2 \left( (A^1_\mu)^2 + (A^2_\mu)^2 \right) + m_1 \left[ A^1_\mu \partial_\mu \phi_2 - A^2_\mu \partial_\mu \phi_1 \right] + g_4 m_1 \left\{ \phi_3 \left[ (A^1_\mu)^2 + (A^2_\mu)^2 \right] - A^2_\mu \phi_1 A^1_\mu + \phi_2 A^2_\mu \right\} - \frac{m_2^2}{2} (\phi_3)^2 - \frac{m_2^2 g_4}{2m_1} \phi_3 (\phi_a)^2 - \frac{m_2^2 g_4}{8m_1} (\phi_a \phi_3)^2, \tag{2}$$

where $N^i_c$ is the spin parameter, $\psi$ Fermi field of the hole, $m_1 = \mu g_4$, $m_2 = 2\sqrt{2} \lambda \mu$, and $T_a$ are the SU(2) generators. The effective Lagrangian describes two massive vector field $A^1_\mu$ and $A^2_\mu$, and one massless U(1) gauge field $A^3_\mu$. Although we have explicitly broken isotopic invariance, the effective Lagrangian density is invariant under local gauge transformations, $\omega(r) = 1 + \alpha^a(r) T_a$, tending to unity at infinity. We shall give the explicit form of the gauge transformations in new variables, confining ourselves to infinitesimal transformations, $\delta \phi_a = -g_4 \varepsilon_{abc} \phi_b \alpha^c - m_1 \varepsilon_{a3c} \alpha^c$.

The generation function $Z[J]$ for Green functions is shown as follows,

$$Z[J] = \int DADBNDCD\tilde{C}D\psi^+ D\psi D\phi \cdot \exp i \int d^4 x \left( L_{\text{eff}} + L_{\text{GF+FP}} + J \cdot \Phi \right). \tag{3}$$

$$L_{\text{GF+FP}} = B^a \partial^\mu A^a_\mu + \frac{1}{2} \alpha B^a B^a + i \tilde{C}^a \partial^\mu D_\mu C^a, \tag{4}$$

where $B^a$ and $C^a$ are Nakanishi-Lautrup(NL) fields and Faddeev-Popov fictitious fields, respectively.

$$J \cdot \Phi \equiv J^\mu_\mu A^a_\mu + J^\mu_\mu B^a + J^\mu_N N_a + \tilde{J}^\mu_C : C^a + J^\mu_\tilde{C} \tilde{C}^a + \tilde{\eta} \psi + \eta \psi^+ + J^\mu_0 \phi_a \tag{5}$$

BRS-quartet [25, 26] in the present theoretical system are $(\phi_1, B^1, C^1, \tilde{C}^1)$, $(\phi_2, B^2, C^2, \tilde{C}^2)$, and $(A^3_\mu, B^3, C^3, \tilde{C}^3)$. Where $A^a_{\mu \neq 0}$ is the longitudinal component of $A^a_\mu$. So we need these fields for the unification condition, although these fields are unobservable and fictitious ones. It is fascinating that a part of the present theoretical formula is similar to so-called Georgi-Glashow model for unifying weak and electromagnetic interactions [27]. Because masses of $A^1_\mu$ and $A^2_\mu$ are formed through the Higgs mechanism by introducing the hole, the fields $A^1_\mu$ and $A^2_\mu$ exist around the hole within the length of $\sim 1/m_1 \equiv R_c$. The quantized gauge fields $A^a_\mu$ are expressed as

$$A^a_\mu = (2\pi)^{-3/2} \int \left[ a^a(p) e^a_\mu(p) \exp (ipr) + a^{a+}(p) e^a_\mu(p) \exp (-ipr) \right] d^3p/\sqrt{2\omega^a_p},$$
where \( \omega_p^a = \sqrt{p^2 + m_i^2} \), \((a = 1, 2)\) and \( \omega_p^a = \sqrt{p^2} \), \((a = 3)\) are the creation and annihilation operators of the gauge particle \( A_\mu^a \) with momentum \( p \), respectively, and \( c_p^a(p) \) are the polarization vectors. The masses, \( m_1 \), of the gauge fields \( A_\mu^1 \) and \( A_\mu^2 \) induced by the hole depend strongly on an angle of Fermi momentum of the hole on the Fermi surface. The value, \( m_1 \), is higher in the case of the hole around the hot spot. It is thought that the interaction between two holes for the Cooper pair formation is derived from the exchange of the fields \( A_\mu^a \) [4-10]. The effective interaction between two holes at \( k \sim (- | \hat{k}_p^h |, 0) \) and \( k \sim (| \hat{k}_p^h |, 0) \) for the Cooper pair formation is given approximately as

\[
H_{\text{int}} \sim g_2^2 \frac{\left( \frac{\omega_{2k_p^h}}{\epsilon_{k_p^h} - \epsilon_{-k_p^h}} \right)^2 - \left( \frac{\omega_{2k_p^h}}{\epsilon_{k_p^h} - \epsilon_{-k_p^h}} \right)}{2} C_{k_p^h}^+ C_{-k_p^h}^+ C_{-k_p^h} C_{k_p^h}.
\]

Where \( \hat{k}_p^h \) is the Fermi momentum at the hot spot, and \( C_{k_p^h}^+ \) and \( C_{k_p^h} \) are the creation and annihilation operators of the hole with momentum \( \hat{k}_p^h \), respectively.

\[
\omega_{2k_p^h} = \sqrt{(2k_p^h)^2 + m_1^2} = \sqrt{(2k_p^h)^2 + \left( \mu(\hat{k}_p^h)g_4 \right)^2}.
\]

The value of \( m_1 \) becomes higher around the hot spot. This means that the attractive interactions between two holes at \( k \sim (- | \hat{k}_p^h |, 0) \) and \( k \sim (| \hat{k}_p^h |, 0) \) or between two holes at \( k \sim (0, - | \hat{k}_p^h |) \) and \( k \sim (0, | \hat{k}_p^h |) \) for the Cooper pair formation are the strongest ones in comparison with the other pairs. As described previously [6,7], the effective interaction between two holes is like the asymptotic freedom, since the SU(2) Yang-Mills fields is more effective when the distance between the holes is shorter than \( \sim 2/m_1 \). For simplifying discussions, we parametrize \( m_1(k_F) = \mu(k_F) \cdot g_4 \) by \( k_F(\theta) \) and angle \( \theta \), with \( \theta = 0 \) corresponding to \( \hat{k} \) along the zone diagonal and represents approximately as \( m_1(k_F) = \mu(k_F) \cdot g_4 = \mu(k_F) \cdot g_4 \cdot \sin^2(2\theta) \). Where \( k_F \) is the Fermi momentum at the hot spot.

3. Restoration of the spontaneous symmetry breaking

The thermodynamic potential is derived from the partition function, That is, \( \Omega(T, \phi_a \phi_a) = -T \ln Z/V \). In the mean field and high temperature approximations, the thermodynamic potential can be introduced as follows,

\[
\Omega(T, \phi_a \phi_a) \sim -2\lambda^2 \mu \phi_a \phi_a + \lambda^2 (\phi_a \phi_a)^2
- P_1(T, m_1) - P_2(T, m_1)
- P_3(T, 0) - P_\phi(T, m_2)
\]

Where \( P_1(T, m_1), P_2(T, m_1), P_3(T, 0), \) and \( P_\phi(T, m_2) \) are the thermal pressures for the massive fields \( A_\mu^1, A_\mu^2 \), the massless field \( A_\mu^3 \), and the massive field \( \phi_3 \), respectively. The thermal pressures for three polarization degrees of the massive gauge fields \( A_\mu^1 \) and \( A_\mu^2 \) are represented approximately as

\[
P_1(T, m_1) = P_2(T, m_1)
\sim \frac{\pi^2}{30} T^4 - \frac{1}{8} m_1^2 \cdot T^2.
\]

The thermal pressure for two polarization degrees of the massless gauge field \( A_\mu^3 \) is represented approximately as

\[
P_3(T, 0) \sim \frac{\pi^2}{45} T^4.
\]
The thermal pressure for one degree of the massive field \( \phi_3 \) is represented approximately as

\[
P_{\phi}(T, m_2) \sim \frac{\pi^2}{90} T^4 - \frac{1}{24} m_2^2 \cdot T^2.
\]

From eq.(7), (8), and (9), the thermodynamic potential \( \Omega(T, \phi_a \phi_a) \) is given by

\[
\Omega(T, \phi_a \phi_a) = \lambda^2 (\phi_a \phi_a)^2 + (-2 \lambda \mu (k_F)^2)
+ \frac{1}{4} \cdot g_3^2 \cdot T^2
+ \frac{1}{3} \lambda^2 \cdot T^2 \phi_a \phi_a - \frac{1}{10} \pi^2 \cdot T^4
\]

(10)

From eq.(10), it is seen that the thermodynamic potential is a function of \( \phi_a \phi_a \) and temperature \( T \). The scalar value, \( \mu(k_F) \), is dependent on the Fermi momentum vector \( \hat{k}_F \). As temperature increases, the minimum of the thermodynamic potential shifts to smaller values of \( \phi_a \phi_a \), and the minimum becomes less deep. The location of the minimum is

\[
(\phi_a \phi_a)_{\text{min}} \equiv \mu(\hat{k}_F, T)_{\text{eff}}^2
= \mu(\hat{k}_F)^2 - \left( \frac{g_2^2}{8\lambda^2} + \frac{1}{6} \right) \cdot T^2
\]

(11)

From eq.(11), the effective mass \( m_{1, \text{eff}}(\hat{k}_F, T) \) of \( A_{1 \mu} \) and \( A_{2 \mu} \) is introduced as

\[
m_{1, \text{eff}}(\hat{k}_F, T) = g_4 \cdot \mu(\hat{k}_F, T)_{\text{eff}}^2
= g_4 \cdot [\mu(\hat{k}_F)^2 \sin^4(2\theta)]
- \left( \frac{g_2^2}{8\lambda^2} + \frac{1}{6} \right) \cdot T^2
\]

(12)

When temperature \( T \) increases above the restoration temperature \( T_{\text{res}}(\hat{k}_F) \equiv \mu(\hat{k}_F) \cdot \left( \frac{g_4^2}{8\lambda^2} + \frac{1}{6} \right)^{-1/2} \approx \mu(\hat{k}_F) \cdot \sin^2(2\theta) \cdot \left( \frac{g_2^2}{8\lambda^2} + \frac{1}{6} \right)^{-1/2} \), the spontaneously broken symmetry is restored. As a result, the effective mass, \( m_{1, \text{eff}}(\hat{k}_F, T) \), of \( A_{1 \mu} \) and \( A_{2 \mu} \) annihilates.

4. An Universal picture of anomalous properties in high-Tc cuprates

In this section, we shall give an universal picture of anomalous properties such as anomalous transport properties, anomalous optical-properties, and anomalous magnetic-properties, from the standpoint of the effect of the quantized massive gauge fields around doped holes and restoration of the spontaneous symmetry breaking in high-Tc cuprates. First, we shall consider the anomalous transport properties in high-Tc cuprates. We use the Keldysh formula [28, 29] for the inequilibrium process in which one of the relaxations of the hole Green function is due to the emission and absorption of the massive gauge fields \( A_{1 \mu} \) and \( A_{2 \mu} \), whose mass \( m_{1, \text{eff}}(\hat{k}_F) \sim g_4 \cdot \mu(\hat{k}_F, T)_{\text{eff}} \). The retarded and advanced in-plane hole Green function \( G_{R,A}(p, \varepsilon) \) is \( 1/(\varepsilon - \xi_p \pm \Sigma) \), where \( \xi_p = (p^2 - k_F^2)/2m^* \), \( \Sigma \) is the self-energy, \( k_F \) is the Fermi momentum, and \( m^* \) is the effective mass of the hole. From Eq.(3), we can obtain the Green function of the massive gauge fields \( A_{1 \mu} \) and \( A_{2 \mu} \) using 't Hooft-Feynman gauge as follows, that is, the Fourier
transform of \( \langle A_{\mu}^a, A_{\nu}^b \rangle_{a=1,2} \) is \( D_{R(k,\omega)} \sim g_{\mu\nu}/(\omega^2 - (k^2 + m_1^2) + \Pi) \). According to the Keldysh formula [30-32], the hole energy relaxation time \( \tau_e \) is defined by the following kinetic equation

\[
\frac{1}{\tau_e} = -\frac{\delta}{\delta n_e} \int dp \frac{1}{(2\pi)^3} \text{Im} \left[ G^A(p, \varepsilon) \right] I(p, \varepsilon),
\]

where \( n_e \) is the hole energy distribution function. In random-phase approximation, the collision integral \( I(p, \varepsilon) \) is represented as

\[
I(p, \varepsilon) = -2 \int dq dq^* \frac{1}{(2\pi)^4} \text{Im} \left[ G^A(p + q, \varepsilon + \omega) \right] \cdot g_2^2 \cdot \text{Im} \left[ D_{R(q,\omega)} R_T(p, q, \varepsilon, \omega) \right].
\]

\[
R_T = \left[ 1 + N_T(\omega) \right] f(p + q, \varepsilon + \omega) \left[ 1 - f(p, \omega) \right] - N_T(\omega) \left[ 1 - f(p + q, \varepsilon + \omega) \right] f(p, \omega),
\]

where \( N_T(\omega) = -\frac{1}{2} \left[ 1 - \coth \frac{\omega}{2T} \right] \) and \( f \) is the hole distribution function. Then the collision integral \( I(p, \varepsilon) \) equals

\[
I(p, \omega) = -2 \int dq dq^* \frac{1}{(2\pi)^4} \text{Im} \left[ G^A(p + q, \varepsilon + \omega) \right] \cdot g_2^2 \cdot \frac{\text{Im} \Pi}{(\omega^2 - (q^2 + m_1^2) + \text{Re} \Pi)^2 + (\text{Im} \Pi)^2} \cdot R_T(p, q, \varepsilon, \omega).
\]

Since the important energy region of a hole is at approximately the Fermi energy, we can consider the collision integral under the condition of \( \varepsilon \sim \omega \ll \sqrt{q^2 + m_1^2} \). In addition, taking account of \( \text{Im} \Pi \sim m_1 < \sqrt{q^2 + m_1^2} \), the collision integral \( I(p, \varepsilon) \) is represented as

\[
I(p, \varepsilon) \sim -2 \int dq dq^* \frac{1}{(2\pi)^4} \text{Im} \left[ G^A(p + q, \varepsilon + \omega) \right] \cdot g_2^2 \cdot \text{Im} \left[ D_{R(q,\omega)} R_T(p, q, \varepsilon, \omega) \right] \cdot m_1(p) g_2^2 \int dq dq^* \text{Im} \left[ G^A(p + q, \varepsilon + \omega) \right] \cdot R_T(p, q, \varepsilon, \omega).
\]

Now we shall show the collision integral \( I(\hat{k}_F, \varepsilon) \) of the Fermi momentum \( \hat{k}_F \). Then, assuming that \( \theta \)-dependence of the mass, \( m_1(\hat{k}_F) \), of the massive gauge fields \( A_{\mu}^1 \) and \( A_{\mu}^2 \) is \( \mu(\hat{k}_F) \cdot g_4 \cdot \sin^2(2\theta) \), the collision integral \( I(\hat{k}_F, \varepsilon) \) due to the massive gauge fields is calculated approximately as

\[
I(\hat{k}_F, \varepsilon) \sim m_1(\hat{k}_F) g_2^2 \int dq dq^* \text{Im} \left[ G^A(\hat{k}_F + q, \varepsilon + \omega) \right] \cdot R_T(\hat{k}_F, q, \varepsilon, \omega) \sim g_4 \cdot \left[ \mu(\hat{k}_F) \right]^2 - \left( \frac{g_4^2}{8\lambda^2} + \frac{1}{6} \right) \cdot T^2 \cdot \frac{1}{2} \cdot g_2^2 \int dq dq^* \text{Im} \left[ G^A(\hat{k}_F + q, \varepsilon + \omega) \right] \cdot R_T(\hat{k}_F, q, \varepsilon, \omega).
\]
\[ R_T(\hat{k}_F, q, \varepsilon, \omega) \]
\[ \sim g_4 \cdot \left[ \mu(\hat{k}_F)^2 \cdot \sin^4(2\theta) \right] - \left( \frac{g_4^2}{8\lambda^2} + \frac{1}{6} \right) \cdot T^2 \right]^{1/2} \cdot g_2^2 \int_{0}^{\infty} dq \omega \text{Im} \left[ G^A(\hat{k}_F + q, \varepsilon + \omega) \right] \cdot R_T(\hat{k}_F, q, \varepsilon, \omega) \] (18)

Since the value, \( \mu(\hat{k}_F) \), is much higher around the hot spot, the value of the collision integral \( I(\hat{k}_F, \varepsilon) \) for the hole of the Fermi momentum \( \hat{k}_F \) near \((\pi, 0)\) or \((0, \pi)\) becomes higher. This implies that the hole lifetime near the \((\pi, 0)\) or \((0, \pi)\) is unusually short. On the other hand, the value, \( \mu(\hat{k}_F) \), is much reduced around the cold spot. As a result, the collision integral \( I(\hat{k}_F, \varepsilon) \) for the hole of the Fermi momentum \( \hat{k}_F \) parallel to \((\pi, \pi)\) decreases remarkably. This means that the hole lifetime in cold spots near the zone diagonal is much longer than elsewhere on the Fermi surface. When temperature \( T \) increases above the restoration temperature \( T_{res}(\hat{k}_F) \equiv \mu(\hat{k}_F) \cdot \left( \frac{g_4^2}{8\lambda^2} + \frac{1}{6} \right)^{-1/2} - \mu(\hat{k}_F) \cdot \sin^2(2\theta) \cdot \left( \frac{g_4^2}{8\lambda^2} + \frac{1}{6} \right)^{-1/2}, \) the spontaneously broken symmetry is restored. As a result, the effective mass, \( m_{1,eff}(\hat{k}_F, T) \), of \( A_1^p \) and \( A_2^p \) annihilates and the anomalous collision integral annihilates above the restoration temperature \( T_{res}(\hat{k}_F) \). That is, the restoration temperature \( T_{res}(\hat{k}_F) \) increases gradually in the Fermi momentum from the cold spot \((\pi, \pi)\), whose \( \theta = 0 \), to the hot spot \((\pi, 0)\), whose \( \theta = 45^\circ \). Thus it is thought that the restoration temperature \( T_{res}(\hat{k}_F) \) around the hot spot \((\pi, 0)\) corresponds to the temperature of large pseudogap (LPG). On the other hand, it is suspected that incoherent Cooper pairs begin to be created around the hot spot, whose angle \( \theta = 45^\circ \) just below the temperature \( T^* \) of small pseudogap (SPG). Then formation of incoherent Cooper pairs extends gradually into the lower angle as temperature decreases from \( T^* \) to \( T_c \). This might correspond to the reduction of the Fermi arc as temperature decreases from \( T^* \) to \( T_c \). We can explain the temperature-evolution of the area of metallic-like Fermi surface naturally from the standpoint of the restoration of the spontaneous symmetry breaking. When \( \mu(\hat{k}_F) \times T_{res}(\hat{k}_F) \) increases proportionally from \( \theta = 0 \) to \( \theta = 45^\circ \) approximately, the evolution of the area of metallic-like Fermi surface is approximately temperature-linear in the range from \( \theta = 0 \) to \( \theta = 45^\circ \). In this condition, it is understood qualitatively that the electrical resistivity becomes approximately linear in temperature, because the quasiparticle lifetime on the area of metallic-like Fermi surface is taken to be \( T^{-2} \) dependence \[17\].

Now, we shall consider anomalous optical-properties in high-Tc cuprates. Hwang et.al.[21] have found the broad background of the optical self-energy spectrum by Fourier-transform infrared (FTIR) spectroscopy is a universal property of the copper-oxygen plane. The optical conductivity is determined by Kramers-Kronig analysis. The complex optical conductivity can be written as \( \sigma(\varepsilon) \sim -i\omega_p^2/[4\pi(2\Sigma^{op}(\varepsilon) - \varepsilon)] \), where \( \omega_p \) is the plasma frequency, \( \Sigma^{op}(\varepsilon) \) is the optical single-particle self-energy, and \( \varepsilon \) is the optical frequency. Imaginary part, \( \text{Im}\Sigma^{op}(\varepsilon) \), of the optical self-energy is \( \sim -1/[2\tau(\varepsilon)] \), where \( \tau \) is the hole energy relaxation time. Using the massive gauge fields \( A_1^p \) and \( A_2^p \) quantized in the path-integral method \[8-10\], the hole energy relaxation time \( \tau_\varepsilon \) is defined by the following kinetic equation,
\[
\frac{1}{\tau_\varepsilon} \propto -\frac{\delta}{\delta n_\varepsilon} \int d\hat{k}_F \left( \frac{1}{(2\pi)^3} \text{Im} \left[ G^A(\hat{k}_F, \varepsilon) \right] I(\hat{k}_F, \varepsilon) \right),
\]
represented as follows,

\[ I(\hat{k}_F, \varepsilon) \propto m_1(\hat{k}_F) g_2^2 \cdot \int dq d\omega \text{Im} \left[ G^A(\hat{k}_F + q, \varepsilon + \omega) \right]. \]

\[ R_T(\hat{k}_F, q, \varepsilon, \omega) = \frac{[1 + N_T(\omega)]f(p + q, \varepsilon + \omega)[1 - f(p, \omega)]}{N_T(\omega)[1 - f(p + q, \varepsilon + \omega)]f(p, \omega)}, \]

where \( N_T(\omega) = -\frac{1}{2} \left[ 1 - \coth \frac{\omega}{2T} \right] \), and \( f \) is the hole distribution function. Since the mass, \( m_1(\hat{k}_F) \), of the massive gauge fields \( A_1^a \) and \( A_2^a \) is relatively large, imaginary part, \( \text{Im} \Sigma^{op}(\varepsilon) \), of the optical self-energy is large. Real part, \( \text{Re} \Sigma^{op}(\varepsilon) \), of the optical self-energy of \( Bi2212 \) shows that the sharp feature (the magnetic resonance peak) can be separated from a broad background and weakens with doping before disappearing completely at a critical doping level of 0.23 holes per copper atom [21]. Superconductivity is still strong in terms of the transition temperature at a critical doping. Thus, Hwang et al. [21] have stressed that both the magnetic resonance peak and phonons are ruled out as the principal cause of high-Tc superconductivity and preferably the broad background is much related to the cause of high-Tc superconductivity. Imaginary part, \( \text{Im} \Sigma^{op}(k, \varepsilon) \), and real part, \( \text{Re} \Sigma^{op}(k, \varepsilon) \), of the optical self-energy are given by,

\[ \text{Im} \Sigma^{op}(k, \varepsilon) \propto -g_2^2 \int dp_1 \int_{-\infty}^{\infty} d\omega \int_{-\infty}^{\infty} d\varepsilon_1 \frac{\text{Im} D_R(k - p_1, \omega) A(p_1, \varepsilon_1)}{\omega + \varepsilon_1 - \varepsilon - i\delta} \cdot \left( \tanh \frac{\varepsilon_1}{2T} + \coth \frac{\omega}{2T} \right). \]  

(19)

\[ \text{Re} \Sigma^{op}(k, \varepsilon) \propto P \int [\text{Im} \Sigma^{op}(k, \varepsilon)/(\varepsilon - \omega')]d\omega'/\pi. \]  

(20)

Where \( D_R(k, \omega) \sim \{g_{\mu}\omega^2 - (k^2 + m_1^2) + \Pi\} \) is the Fourier transform of the Green function, \( \langle A_1^a A_2^a \rangle_{a=1,2} \), of the massive gauge fields \( A_1^a \) and \( A_2^a \) in the 't Hooft-Feynman gauge. \( A(p_1, \varepsilon_1) \) is the hole spectral function. Since the spectrum function, \( \text{Im} D_R(k - p_1, \omega) \), has the broad distribution, it is regarded that \( \text{Re} \Sigma^{op}(\varepsilon) \equiv \int \text{Re} \Sigma^{op}(k, \varepsilon)dk \) has the broad distribution, whose width is \( \propto m_1 \), from eqs (19) and (20). Now we shall consider the mechanism why the intensity of the broad background in \( \text{Re} \Sigma^{op}(\varepsilon) \) increases as the temperature decreases. Recently the present author [3,4] has introduced the restoration temperature \( T_{res}(\hat{k}_F) \equiv \mu(\hat{k}_F) \cdot (g_{\mu}^2/8\Lambda^2 + 1/6)^{-1/2} \). The restoration temperature \( T_{res}(\hat{k}_F) \) increases gradually in the Fermi momentum from the cold spot \((\pi/2, \pi/2)\) to the hot spot \((\pi, 0)\). When the temperature \( T \) decreases below the restoration temperature, the spontaneously symmetry breaking is occurred. As the temperature decreases, the region of the Fermi momentum \( \hat{k}_F \), which contributes to the broad distribution of \( \text{Re} \Sigma^{op}(\hat{k}_F, \varepsilon) \), is increased due to the evolution of destruction of the Fermi surface. Thus, the intensity of the broad background in \( \text{Re} \Sigma^{op}(\varepsilon) \) increases as the temperature decreases. Then, we discuss the linear-like decrease in the spin part \( K_s \) of the Knight shift as temperature decreases from \( T_{mk} \) to \( T^* \), below which the nuclear spin-lattice relaxation rate divided by temperature, \( 1/(T_1 T) \), starts to decrease gradually [22]. The spin part \( K_s \) of the Knight shift is represented approximately as \( K_s \sim \frac{\gamma_0}{4} |u_F(0)|^2 \chi_P \), by using the block function, \( \psi_0(r) = u_k(r) e^{ikr} \), of the electrons. Where \( \chi_P = 2\mu_B^2 N(E_F) \) is the paramagnetic spin susceptibility, and \( \mu_B \) is Bohr magneton. \( N(E_F) \) is the density of state of electrons at the Fermi energy. When the linear-like reduction
of the Fermi arc with decreasing the temperature is assumed, \( N(E_F) \) decreases linearly as the temperature decreases. As a result, the spin part \( K_\theta \) of the Knight shift decreases linearly from \( T_{nK} \) to \( T^* \). In addition, Oda et al. [33] reported that the uniform susceptibility \( \chi(T) \) starts to decrease gradually below \( T_{m\chi} \), whose value is near that of \( T_{nK} \). This suggests strongly that the restoration temperature \( T_{res}(\tilde{k}_F) \) around the hot spots corresponds to \( T_{nK} \) and \( T_{m\chi} \).

5. Conclusion
We have presented an universal picture of anomalous properties in high-Tc cuprates, from the standpoint of the effect of the quantized massive gauge fields around doped holes and restoration of the spontaneous symmetry breaking. It is suggested strongly that quantized massive gauge fields around doped holes, which have been introduced theoretically in the path-integral method by the present author, are mediating Cooper pairing in high-Tc cuprates.

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