Collective modes in strange and isospin asymmetric hadronic matter

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Abstract

We study the propagation of non-strange and strange meson modes in hadronic matter considering both isospin and strangeness mixings induced by quantum fluctuations in the medium. Baryons are described using the Quark Meson Coupling model extended to include interactions of strange quarks. In particular we evaluate the dependence of the meson masses on the baryonic density, the strangeness fraction and the isospin asymmetry of the medium. We have found a considerable admixture of strangeness and isospin in the $\sigma$-mode in the high density regime.

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The study of the meson properties in a hot/dense hadronic medium is at present an active field of research, since it is related to fundamental aspects of the low energy regime of the strong interactions such as the formation of a plasma of quarks and gluons and the restoration of the chiral symmetry. A prominent role is played by the lightest vector mesons, as they are closely connected to the emission rate of dileptons in heavy ion collision experiments. On the other hand the scalar mesons still keep valuable information about the quark structure of hadrons as can be seen, for example, in the dubious composition of the almost degenerate $a_0(980) - f_0(980)$ mesons.

There exist a profuse quantity of worthful studies of the in-medium modification of meson properties, some of them are based on pure quark models such as the Nambu Jona-Lasinio or use descriptions in terms of hadrons only [1,2,3,4,5,6,7]. There were also some efforts to reconcile both aspects by using hybrid models such as the Quark Meson Coupling (QMC) [8].

A very interesting feature appearing in these treatments is the mixing effect, which combines states of different isospin or Lorentz components [4,5,6]. This dynamical violation of the lagrangian invariances is particularly important in considering meson propagation, because it opens new in-medium decaying channels.

The aim of this work is to study qualitatively the modification of the meson properties in a dense medium as described by the baryonic collective modes in the relativistic random phase approximation (RRPA). In order to

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adjust to the conditions found in the experience, as for instance in relativistic heavy ion collisions, we consider matter with diverse strangeness fraction as well as isospin asymmetry. We take into account the quark structure of baryons by using the QMC model [8,9,10]. In our treatment we regard all the mixing polarizations allowed by the hadronic lagrangian proposed.

The evaluation of the RRPA with both isospin and strangeness mixing effects is still lacking at present, and the qualitative description of the lightest hadron masses in such a wide variety of situations as found in the experience is also deficient.

The QMC model relates the in-medium properties of hadrons with their quark structure. Baryons are represented as non-overlapping spherical bags containing three valence quarks, while the bag radius changes dynamically with the medium density. The exchange of mesons coupled directly to the confined quarks provide the baryon-baryon interaction.

To fit the model to our purposes we consider the strange $f_0(980)$ and $\phi$ mesons as in Ref. [10], and the scalar iso-vector $a_0(980)$ together with the more commonly used $\sigma$, $\omega$, and $\rho$ mesons.

The mean field approximation (MFA) equation for a quark of flavor $q$ ($q = u, d, s$), of current mass $m_q$ and $I_3$ third isospin component is

$$
(i\gamma^\mu \partial_\mu - m_q^* - g_\omega^q \gamma^0 \bar{\omega} - g_\phi^q \gamma^0 \bar{\phi} - g_\rho^q I_3^q \gamma^0 \bar{\rho})\Psi_q = 0,
$$

(1)

where we have used

$$
m_q^* = m_q - g_\omega^q \bar{\sigma} - g_\phi^q \bar{\zeta} - g_\rho^q I_3^q \bar{\delta}.
$$

(2)

In these Eqs. $\delta$ and $\zeta$ stand respectively for the $a_0$ and $f_0$ fields, the upper bar indicates mean values. The solution of Eq. (1) for a quark confined within a spherical bag of radius $R_b$, representing a baryon of class $b$, is well known and can be consulted for example in [8,9]. The baryon effective mass $M_b^*$ is given by

$$
M_b^* = \frac{\sum_q N_q^b \Omega_{qb} - z_{0b}}{R_b} + \frac{4}{3} \pi BR_b^3,
$$

(3)

where $N_q^b$ is the number of quarks of flavor $q$ inside the bag and $\Omega_{qb} = [x_{qb}^2 + (R_b m_q^*)^2]^{1/2}$. The eigenvalue $x_{qb}$ comes from the so called linear boundary condition at the bag surface [9]. The bag constant $B^{1/4} = 210.86 MeV$ stands for non-perturbative vacuum contributions, it is adjusted to get a proton bag radius $R_p = 0.6 fm$ in vacuum. The zero-point motion parameters $z_{0b}$ are fixed to reproduce the baryon spectrum at zero density. In our calculations we have taken $m_u = m_d = 5 MeV$ and $m_s = 150 MeV$ for the current quark masses.

The equilibrium condition for the bag radius is [11]

$$
- \frac{1}{4\pi R_b^2} \left( \frac{\partial M_b^*}{\partial R_b} \right) = \frac{1}{3\pi^2} \sum_{q'} k_{q'} \int_0^{k_{q'}} \frac{dk k^4}{\sqrt{M_{q'}^{*^2} + k^2}}
$$

(4)

The sum on the r.h.s runs over all the baryonic species considered.

We neglect the coupling between strange and non-strange quarks and mesons. To determine the remaining couplings we assume pure vector dominance together with $SU(3)$ flavor symmetry, which allows to regain the
The self-consistent equations for the meson mean fields are of the form

\[ U^\alpha = \text{potential depth} = \sum g^\alpha_{\sigma,\omega} \Xi \]

The one-loop proper polarization insertion describing the propagation between meson states \( \sigma, \omega \) is given by

\[ \Pi_{\alpha\beta}(q) = -i \sum_b g^b_{\alpha\beta} \int \frac{d^4k}{(2\pi)^4} \text{Tr} [G_b(q) \Gamma_{\alpha} G_b(q + k) \Gamma_{\beta}] \]

(7)

where \( N^{b}_{n,s} (N^{b}_{s}) \) is the non-strange (strange) quark number inside the baryon \( b \), with third isospin component \( I^b_3 \), and in addition we have \( g^u_{\delta,\rho} = -\sqrt{3} g^u_{\sigma,\omega} \). Furthermore as the \( SU(2) \) flavor in the non-strange sector is almost perfectly realized, we take as identical the coupling of quarks \( u \) and \( d \) with every meson. Therefore we have five independent couplings \( g^u_{\sigma,\omega}, g^u_{\rho,\phi}, g^d_{\sigma,\omega}, g^d_{\rho,\phi}, \) and \( g^s_{\sigma,\omega} \), we fix their numerical values by fitting symmetric nuclear matter properties at saturation, namely baryonic density \( n_0 = 0.15 \text{fm}^{-3} \), binding energy \( E_b = -16 \text{MeV} \) and symmetry energy \( a_s = 35 \text{MeV} \). The last condition gives a functional relationship between \( g^u_{\delta,\rho} \) and \( g^d_{\rho,\phi} \), we have chosen an arbitrary point on this curve. To fix \( g^s_{\sigma,\omega} \) we follow references [3,10,14] assuming a potential depth \( U_\Xi = 40 \text{MeV} \) for the \( \Xi \) hyperon embedded in symmetrical \( \Xi \) matter at baryonic density \( n_\Xi = n_0 \). Thus we obtain \( g^u_{\sigma,\omega} = 5.99, g^u_{\rho,\phi} = 3.00, g^d_{\sigma,\omega} = 1.83, g^d_{\rho,\phi} = 4.50, \) and \( g^s_{\sigma,\omega} = 4.48 \).

The self-consistent equations for the meson mean fields are of the form

\[ \bar{\chi} = -\frac{1}{m^2 \chi} \sum_b Y^b_{\chi} n^b_{\chi} \]

(6)

with \( Y^b_{\chi} = \partial M^*_b / \partial \bar{\chi}, n^b_{\chi} = n^b_{\chi} \) if the generic meson \( \chi \) is a Lorentz scalar field, or \( Y^b_{\chi} = -g^u_{\sigma,\omega}, -g^d_{\sigma,\omega}, -g^d_{\rho,\phi}, n^b_{\chi} = n^b_{\chi} \) for the vector case. The scalar and baryon densities

\[ n^b_{\sigma} = \frac{1}{\pi^2 M^*_b} \int_0^{k_{F_b}} \frac{dk}{\sqrt{M^*_b}^2 + k^2}, \quad n^b_{\rho,\phi} = \frac{k_{F_b}^3}{3\pi^2} \]

are functions of the Fermi momentum \( k_{F_b} \). The meson quantum fluctuations will be considered later. For the hadronic masses in vacuum we have used the values \( m_\sigma = 550, m_\omega = 980, m_\rho = 783, m_\phi = 1020, m_\delta = 984, m_\rho = 770, M_\rho = M_\pi = 938.92, M_\Lambda = 1115.63, \) and \( M_\Xi = 1318.11 \), expressed in MeV.

The QMC model describes effective baryons propagating in a homogeneous background of classical meson fields as described above. Beyond this picture we assume a linear coupling between the mesonic fluctuations and the baryons emerging from the MFA, with coupling constants given in Eq. (5). The full meson propagator can be approximated in a non-perturbative approach by summing the ring diagrams to all orders, i.e. the RRPA. This is a common procedure in nuclear physics, and can be used to evaluate the density dependence of the meson masses among other properties. The QMC has been used previously to evaluate the vector meson masses, regarding bosons as bags on the same foot as the baryonic ones, but neglecting the quark structure of the \( \sigma \) meson [8].

The one-loop proper polarization insertion describing the propagation between meson states \( \alpha \) and \( \beta \) is given by

\[ \Pi_{\alpha\beta}(q) = -i \sum_b g^b_{\alpha\beta} \int \frac{d^4k}{(2\pi)^4} \text{Tr} [G_b(q) \Gamma_{\alpha} G_b(q + k) \Gamma_{\beta}] \]

(7)
with indexes $\alpha, \beta$ running over the whole set of mesons and its internal degrees of freedom. Further, $G_b(p)$ stands for the baryon $b$ propagator in the MFA, and the structure of $\Gamma_\alpha$ depends on the baryon-meson vertex, namely $\Gamma_{\sigma, \zeta} = 1, \Gamma_\delta = \tau, \Gamma_{\omega, \phi} = \gamma^\mu, \Gamma_\rho = \gamma^\mu \tau$. Eq. (7) contains ultraviolet divergences coming from the vacuum contribution which require an appropriate regularization. We adopt the dimensional regularization procedure, preserving that $\Pi_{\alpha \beta}(q) = \Pi_{\beta \alpha}(q)$, and we shall discuss later the choice of the regularization points. Explicit expressions for Eq. (7) are similar to that given in [2,6], for instance.

The formalism is best described within a generalized meson propagator in a matrix representation of dimension equal to the sum of the mesonic degrees of freedom. For example the free ($P^0$), and full ($P$) generalized meson propagators have respectively in its diagonal blocks the free and dressed meson propagators of the $\sigma, \zeta, \delta, \omega, \phi$, and $\rho$ fields, and in the complementary spaces they have zeros or finite mixing polarizations arising from Eq. (7), respectively. The corresponding Dyson-Schwinger equation in matrix form can be used to solve for $P(q)$:

$$P(q) = [I - P^0(q) \Pi(q)]^{-1} P^0(q)$$

(8)

The dielectric function $\epsilon(q) = \text{det} [I - P^0(q) \Pi(q)]$ is defined so that its roots coincide with the poles of $P(q)$. We study the temporal regime $\vec{q} = 0$, where scalar and vector contributions to the equation $\epsilon(q) = 0$ become independent; in addition the longitudinal and transversal branches of the vector mesons become degenerate. Hence we have two separate equations of the form:

$$(D^{-1}_\lambda - \Pi_{\lambda \lambda})(D^{-1}_\mu - \Pi_{\mu \mu})(D^{-1}_\nu - \Pi_{\nu \nu}) - (D^{-1}_\lambda - \Pi_{\lambda \lambda})\Pi_{\mu \nu} - (D^{-1}_\nu - \Pi_{\nu \nu})\Pi_{\lambda \mu} = 0,$$

(9)

where the indices $\lambda, \mu, \nu$ take definite meson labels, namely $\lambda = \delta, \mu = \sigma$, and $\nu = \zeta$ for the scalar case, and $\lambda = \rho, \mu = \omega$, and $\nu = \phi$ for the vector case.

Furthermore $D^{-1}_\alpha = q^2 - m^2_\alpha$ is the inverse free propagator for the $\alpha$-meson. It must be noted that there is no $\delta-\zeta$ neither $\rho-\phi$ mixing, and this causes the splitting of the isovector mode in vacuum as well as in the isospin symmetric dense medium. This splitting appears because the $\sigma-\delta$ and $\omega-\phi$ polarizations become null under these conditions, in view of the degeneracy assumed for the baryonic iso-multiplets.

In order to extract finite values from the polarization we require at zero baryonic density [12]

$$\Pi_{\alpha \beta}(q^2 = R^2_{\alpha \beta}) = 0,$$

$$\left(\partial \Pi_{\alpha \beta}(q)/\partial q^2\right)_{q^2=R^2_{\alpha \beta}} = (1 - L_{\alpha \beta}) \sum_b g^{b}_{\alpha} g^{b}_{\beta}/(8\pi^2)$$

The last equation holds only for both $\alpha$ and $\beta$ corresponding to scalar mesons, where $g^{b}_{\alpha} = g^{b}_{\alpha}$ for the isoscalars $\alpha = \sigma, \zeta$, or $g^{b}_{\alpha} = g^{b}_{\alpha} P^{1}_{\lambda}$ for the isovector $\alpha = \delta$. The regularizing parameter is fixed at $L_{\alpha \beta} = 10$ [7,12]. Here $R^2_{\alpha \beta} = R_{\alpha \beta}$ are a set of regularization points that we choose in order to reproduce the physical meson masses at zero density. This requirement determines unambiguously $R_{\delta \delta} = m_{\delta}$, and $R_{\rho \rho} = m_{\rho}$, otherwise we obtain a range of points $(R_{\sigma \sigma}, R_{\sigma \zeta}, R_{\zeta \zeta})$ and $(R_{\omega \omega}, R_{\omega \phi}, R_{\phi \phi})$. The regularization points for the mixing polarizations including one iso-vector meson are not constrained, for the reasons mentioned in the preceding paragraph. In our calculations we have used $R_{\sigma \sigma} = m_{\sigma}, R_{\sigma \zeta} = R_{\zeta \delta} = 2.807, R_{\zeta \zeta} = 4.654, R_{\omega \omega} = 3.98, R_{\omega \phi} = R_{\omega \rho} = 6.09, R_{\phi \phi} = 5.12$ expressed in $fm^{-1}$. We have numerically checked that the final conclusions are insensitive to the precise choice of these parameters. Instead, the choice of $L$ has noticeable effects on the behavior of the $\sigma$ meson properties [7,12], and a detailed discussion within this context will be given elsewhere [13].
To study the behavior of meson masses in a hadronic medium with strange and isospin content we have considered two situations: a) nucleons and Λ, b) nucleons, Λ and Ξ in equilibrium against strong decay. In the last case the relative abundance of Λ and Ξ is determined by the relation between their chemical potentials, a situation often treated in the literature [3,10,14]. These two different hadronic environments when considered at zero isospin asymmetry give practically identical results for the meson masses. Therefore we consider in the following only case (a).

We take the total baryonic density $n = (n^n + n^p + n^\Lambda)$ running up to $4n_0$, guarantying that baryonic bags do not overlap. For a fixed total baryonic density $n$ we introduce the isospin asymmetry parameter $t = (n^n - n^p)/n$ and the strangeness fraction $S = n^\Lambda/n$, both $S$ and $t$ taking values in $(0, 1)$.

In Fig. 1 we plot the effective meson masses computed as the zeros of Eq.(9), for $t = 0$ and different strangeness fractions $S$. Although the meson quantum numbers are blurred in-medium due to the mixing effect, we can distinguish well defined masses branches continuously related to its zero density value. Therefore we keep the label of the meson which originates each branch.

To explain the behavior of the meson masses shown in this figure one must take into account that the polarization insertion is composed of two additive terms of opposite sign, namely the vacuum contribution $\Pi^{(F)}$ and the density dependent part $\Pi^{(D)}$ which explicitly depends on the Fermi momenta of the baryons. Thus, neglecting mixing corrections one has $m_k^s = m_k^0 + \Pi^{(F)}_k + \Pi^{(D)}_k$, $k = \zeta, \delta, \omega, \phi$ and $\rho$. At low densities ($n \leq n_0$) vacuum fluctuations are dominant and tend to lessen the meson effective masses. However the density dependent part of the polarizations contributes in the opposite direction, attenuating or even reverting this trend at higher densities. The only exception is given by $m_{\sigma}^s$ for which $\Pi^{(F)}$ raises at low densities due to the relatively small value of $m_{\sigma}$.

Within this approach the dependence on $S$ for a fixed baryonic density can be explained by observing that $\Pi^{(F)}$ is a negative decreasing function of $S$. On the other hand $\Pi^{(D)}$ decreases (increases) when $S$ increases for the non-strange (strange) mesons, due to the depletion (population) of the baryons to which they couple. Furthermore the variation of $\Pi^{(D)}$ is stronger for the vector than for the scalar mesons. For instance, as $S$ grows the total polarization of the $\phi$ branch increases due to the predominance of $\Pi^{(D)}_{\phi\phi}$ over $\Pi^{(F)}_{\phi\phi}$, therefore $m_{\phi}^s$ is enhanced. On the other hand for the $\zeta$ branch, the growing of $\Pi^{(D)}$ is much weaker and $\Pi^{(F)}$ determines the decreasing behavior of $m_{\zeta}^s$ with $S$. The numerical results shown below provide the justification for disregarding in first approach the mixing effects for the strange mesons. This reasoning does not hold for the $\sigma$ branch because it carries a strong mixing, leading to a more involved behavior, but the final result is a decreasing dependence with $S$.

On the other side, as can be appreciated from this figure, the $\omega - \rho$ mass difference for baryon densities around $2n_0$ is close to the pion mass, increasing the probability of the $\rho \to \omega + \pi$ decay. This fact in turn contributes to enhance the $\rho$ meson width, in agreement with several theoretical predictions as well as with phenomenological observations.

We can compare Fig. 1 with the results shown in Ref. [3] (Figs. 3 and 4 for $q = 1 MeV$) for the $\sigma, \omega$ and $\zeta$ mesons. It can be appreciated that the overall behavior with $S$ agrees with our calculations within the considered range of densities.

The influence of the isospin asymmetry $t$ over the meson modes at constant strangeness $S = 0.5$ is shown in Fig. 2. The qualitative description does not change substantially for other values of $S$. The strange meson masses are almost independent of $t$, as they are coupled only to the iso-singlet $\Lambda$. On the left hand panel, the $\sigma$-meson mass shows a slight enhancement with growing $t$, this effect turns to be appreciable for $n > 2n_0$. On the opposite panel of this figure, the isovector $\rho$-meson mass also exhibits a small enhancement with $t$ for $n > 1.5n_0$. Interestingly the isoscalar $\omega$-meson mass presents the inverse trend when $t$ increases, therefore we conclude that for a fixed strangeness content, the splitting between the $\omega$ and $\rho$ masses increases with the isospin asymmetry. These results can be compared, for example, with those of Ref. [4] where the same behavior of $m_{\omega,\rho}^s$ with $t$ is found.
A measure of the mixing between two physical meson states $\alpha$ and $\beta$ can be given by the mixing angle $\theta_{\alpha\beta}$, defined by [5]

$$\tan 2\theta_{\alpha\beta} = \frac{2\Pi_{\alpha\beta}}{m_\alpha^2 + \Pi_{\alpha\alpha} - m_\beta^2 - \Pi_{\beta\beta}}.$$  \hspace{1cm} (10)

Since our aim is to obtain a qualitative estimation of how much isospin and strangeness remain good quantum numbers for the collective modes, we do not include the width of the original mesons in Eq. (10) in consistency with the RRPA. However quantitative calculations must include an imaginary contribution to the polarizations [4,15].

For $t = 0$ the isoscalar-isovector mixing vanishes, remaining the isoscalar strange-non-strange mixing only. We have evaluated Eq. (10) for $\alpha = \sigma$ $(\omega)$, $\beta = \zeta$ $(\phi)$ at the pole $p_0 = m_\alpha^*$ $(m_\beta^*)$ in terms of the baryonic density for various strangeness fractions $S$. The result is plotted on the upper left (upper right) panel of Fig. 3; as expected the mixing effect increases with $n$ and $S$. We can observe that $\theta_{\sigma\zeta}(q_0 = m_\sigma^*)$ becomes appreciable at $n \gtrsim n_0$, and therefore the mixing effects in the $\sigma$ channel may not be neglected even at relatively low densities. In the case of $\theta_{\omega\phi}(q_0 = m_\zeta^*)$ it turns to be significative only for high densities $(n \gtrsim 3n_0)$. With respect to $\theta_{\sigma\zeta}(q_0 = m_\zeta^*)$ and $\theta_{\omega\phi}(q_0 = m_\phi^*)$ (not shown here), their amplitudes remain below $3.5^\circ$ for all the densities studied.

Because the strangeness mixing remains small for all but the $\sigma$ meson, we can estimate the isoscalar-isovector mixing for $t \neq 0$ by using Eq. (10) also. In the lower left panel of Fig. 3 it can be seen that scalar $\sigma - \delta$ mesons mix only appreciably in the $\sigma$ branch at densities $n \gtrsim 3.5n_0$. Instead, non-strange $\omega - \rho$ vector mesons exhibit a noticeable isospin mixing of approximately equal magnitude in both branches, as shown in the lower right panel of this figure. However, inclusion of the $\omega$ and $\rho$ widths is expected to decrease these values [4,15]. By raising the strangeness fraction a sizeable decrease of the mixing angle is obtained for both scalar and vector mesons, in the considered range of $n$.

We conclude that the in-medium propagation of the $\sigma$-channel concentrates non-negligible admixtures of isospin as well as strangeness in the high density regime. The vector $\omega$ and $\rho$ modes are affected by isospin but not by strangeness mixing. On the other hand the $\delta$, $\zeta$, and $\phi$ modes propagate in almost pure states for all the densities studied.

The effects discussed for vector mesons could be observable for instance, in the low mass dilepton production from heavy ion collision. The in-medium modification of the vector mesons has been successfully invoked to explain the gross features of the dilepton spectra from the CERES collaboration [16], in the frame of the Vector Meson Dominance model. A key role is assigned to the in-medium dropping of the $\rho$-meson mass. In order to get this result at the tree level, a direct coupling of the $\sigma$ to the $\omega/\rho$ fields was proposed [16]. In our approach we obtain the same qualitative behavior within the RPA, and also a description of the mixture of states. Because of the smallness of the $\omega - \phi$ mixing shown in Fig. 3 and the fact that both mesons decay largely after the freeze out, we do not expect a significative modification in these channels. The $\omega - \rho$ mix is large enough at almost all densities to have significative consequences. Because the $\rho$ and $\omega$ mesons decay mainly into two and three pions respectively, in the dense hadronic medium the alternative mechanism $\rho \to \omega \to 3\pi$ is viable, having the effect of removing part of the strength between the peaks corresponding to both mesons and transferring it to the left of the $\rho$ peak in the dilepton spectra.

Some points deserve to be investigated within this model, for example the meson coupling to Goldstone bosons, the fact that the non-overlapping bag hypothesis is violated in the extreme density regime used in astrophysical applications [11], and the medium dependence of the scalar coupling constants which play a significative role in chiral models [17,18].
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Fig. 1. The effective meson masses as a function of the density for $t = 0$ and several strangeness fractions $S = 0, 0.25, 0.5, 0.75, 1$. The left (right) panel corresponds to the scalar (vector) meson branches. The arrows indicate the direction of growing $S$. 
Fig. 2. The effective meson masses as a function of the density for $S = 0.5$ and several isospin compositions $t = 0, 0.5, 1$. The left (right) panel corresponds to the scalar (vector) meson branches. The arrows indicate the direction of growing $t$. 
Fig. 3. The mixing angles defined by Eq. (10) as a function of the density. The upper panels show strangeness mixing for $t = 0$ and several strangeness fractions $S = 0$, 0.25, 0.5, 0.75, 1. The upper left (upper right) panel corresponds to the scalar $\sigma$-$\zeta$ (vector $\omega$-$\phi$) mix evaluated at the $\sigma$ ($\omega$) meson branch. The arrows indicate the direction of growing $S$. Lower panels show the isospin mixing angles for $t = 0.5$ and the strangeness fractions $S = 0$, 0.5. The lower left (lower right) panel corresponds to the scalar $\sigma$-$\delta$ (vector $\omega$-$\rho$) mix evaluated at both the $\sigma$ and $\delta$ ($\omega$ and $\rho$) meson branches.